

Quantum enhanced measurements without entanglement

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Quantum-enhanced measurements exploit quantum mechanical effects for increasing the sensitivity of measurements of certain physical parameters and have great potential for both fundamental science and concrete applications. Most of the research has so far focused on using highly entangled states, which are, however, difficult to produce and to stabilize for a large number of constituents. In the following we review alternative mechanisms, notably the use of more general quantum correlations such as quantum discord, identical particles, or non-trivial hamiltonians; the estimation of thermodynamical parameters or parameters characterizing non-equilibrium states; and the use of quantum phase transitions. We describe both theoretically achievable enhancements and enhanced sensitivities, not primarily based on entanglement, that have already been demonstrated experimentally, and indicate some possible future research directions.

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I. INTRODUCTION

A. Aim and scope

Quantum enhanced measurements (QEM) aim at improving measurements of physical parameters by using quantum effects. The improvement sought is an enhanced sensitivity for a given amount of resources such as mean or maximum energy used, number of probes, number of measurements, and integration time. Ideas in this direction go back at least to the late 1960s when the effect of quantum noise on the estimation of classical parameters started to be studied systematically mathematically (Helstrom, 1969; Holevo, 1982). In the early 1980s

first detailed proposals appeared on how to enhance the sensitivity of gravitational wave detectors by using squeezed light (Caves, 1980, 1981). Nowadays, squeezed light is routinely produced in many labs, and in 2013 squeezed light technology was finally introduced in one of the LIGO gravitational wave observatories (Aasi *et al.*, 2013).

Quantum enhanced measurements have the potential of enabling many important applications, both scientific and technological. Besides gravitational wave detection, there are proposals or demonstrations for the improvement of time- or frequency-standards, navigation, remote sensing, measurement of very small magnetic fields (with applications to medical brain- and heart-imaging), measurement of the parameters of space-time, thermometry, and many more. We do not review the vast literature on quantum metrology in general, as several recent reviews and even monographs are available (Degen *et al.*, 2016; Giovannetti *et al.*, 2006; Paris, 2009; Pezzé and Smerzi, 2014; Pezzé *et al.*, 2016; Tóth and Apellaniz, 2014; Wiseman and Milburn, 2009).

From the theoretical side, the standard tool for evaluating a possible quantum enhancement has become the so-called quantum Cramér-Rao bound (QCRB) (Braunstein and Caves, 1994; Braunstein *et al.*, 1996; Helstrom, 1969; Holevo, 1982). It provides a lower bound on the variance $\text{Var}(\theta_{\text{est}})$ of any unbiased estimator function θ_{est} that maps observed data obtained from arbitrary quantum measurements (in the form of POVM measurements, i.e. measurements even more general than projective von Neumann measurements) to an estimate of the parameter θ . The bound is optimized over all possible POVM measurements and over all unbiased estimator functions, i.e. estimator functions whose average equals the true value of the parameter in at least an infinitesimal interval about the value of θ considered. In the limit of an infinite number of measurements the bound can be saturated. It thus represents a valuable benchmark that can in principle be achieved once all technical noise problems have been solved, such that only the unavoidable noise inherent in the quantum state itself remains. The quantum state is in general a mixed state $\rho(\theta)$, i.e. may contain besides pure quantum noise also uncertainty about the preparation or thermal noise.

The precise form of the QCRB reads

$$\text{Var}(\theta_{\text{est}}) \geq \frac{1}{MI_{\theta}}, \quad (1)$$

with M the number of independent measurements and I_{θ} the so-called quantum Fisher information (QFI). In (Braunstein and Caves, 1994) it was shown that I_{θ} is a geometric measure on how much $\rho(\theta)$ and $\rho(\theta + d\theta)$ differ, where $d\theta$ is an infinitesimal increment of θ . It thus offers the physically intuitive picture that the parameter θ can be measured the more precisely the more strongly the state $\rho(\theta)$ depends on it (see below for a precise def-

inition).

A standard classical method of noise reduction is to average measurement results from N independent, identically prepared systems. Based on (1), one can easily show that this leads to a limitation of smallest possible standard deviation of an unbiased θ_{est} that scales as $1/\sqrt{N}$, known as the so-called standard quantum limit (SQL). In a quantum mechanical formulation with pure states, the situation corresponds to having the N quantum systems in an initial product state, $|\psi\rangle = \otimes_{i=1}^N |\phi\rangle_i$. Suppose that the parameter is encoded in the state through a unitary evolution with a hamiltonian $H(\theta) = \theta \sum_{i=1}^N h_i$, i.e. $|\psi(\theta)\rangle = \exp(-iH(\theta))|\psi\rangle$. Then, with M final measurements, the smallest achievable standard deviation $\sigma(\theta_{\text{est}}) \equiv (\text{Var}(\theta_{\text{est}}))^{1/2}$ of the estimation of θ is

$$\sigma(\theta_{\text{est}})_{\min} = \frac{1}{\sqrt{NM}(\Lambda - \lambda)}, \quad (2)$$

where Λ and λ are the largest and smallest eigenvalue of h_i , respectively, taken for simplicity here as identical for all subsystems (Giovannetti *et al.*, 2006). In fact, this $1/\sqrt{N}$ scaling can be easily understood as a consequence of the central limit theorem in the simplest case that one measures the systems independently. But since (1) is optimized over all measurements of the full system, it also implies that entangling measurements of all systems after the parameter has been encoded in the state cannot improve the $1/\sqrt{N}$ scaling.

Unfortunately, there is no unique definition of the SQL in the literature. Whereas in the described $1/\sqrt{N}$ scaling N refers to the number of distinguishable sub-systems, the term SQL is used for example in quantum optics typically for a scaling as $1/\sqrt{\bar{n}}$ with the average number of photons \bar{n} , which in the same mode are to be considered as indistinguishable (see Sec.III). In this context, the $1/\sqrt{\bar{n}}$ scaling is also called “shot-noise limit”, referring to the quantum noise that arises from the fact that the electromagnetic energy is quantized in units of photons. Furthermore, the prefactor in these scaling behaviors is not fixed. We therefore define here quite generally SQL as the best scaling that can be achieved with all resources in a classical state. While this is not yet a mathematical definition either, it becomes precise once the classical resources are specified in the problem at hand. For example, in quantum optics a pure classical state with \bar{n} photons is a coherent state. Once these resources are specified, it is clear what \bar{n} counts, and also the prefactor of the SQL scaling law becomes fixed. In that manner, quantum enhancements are possible not only by improving the scaling law, but also by changing the prefactor.

One can also show from (1) that initially entangled states can improve the scaling to $1/N$ (Giovannetti *et al.*, 2006), known as the “Heisenberg-limit” (HL). Similarly to the SQL, there is no unique definition of the HL in the literature (see the remarks in sec.IV.B). Nevertheless, achieving “the HL” has been the goal of large

experimental and theoretical efforts over the last two decades. However, only few experiments achieved the $1/N$ scaling of the HL and only for very small numbers of sub-systems, where the scaling advantage is still far from allowing one to beat the best possible classical measurements. This has several reasons: first of all, it is already very difficult to achieve even the SQL, as all non-intrinsic noise sources have to be eliminated. Secondly, resources such as photons are cheap, such that classically one can operate with very large photon numbers, whereas entangled states with large photon numbers are difficult to produce. Thirdly, and most fundamentally, QEM schemes are plagued by decoherence. Indeed, it has been shown that a small amount of Markovian decoherence brings the $1/N$ scaling for certain highly entangled states back to the $1/\sqrt{N}$ scaling of the SQL (Escher *et al.*, 2011; Huelga *et al.*, 1997; Kołodyński and Demkowicz-Dobrzański, 2010). Recent research has focussed on finding optimal states in the presence of decoherence, and at least for non-Markovian noise, a certain improvement can still be obtained from entangled states (Chin *et al.*, 2012; Matsuzaki *et al.*, 2011). Also, niche applications are possible, where the light-intensity must be very small, as in some biological applications. Nevertheless, it appears worthwhile to think about alternative possible quantum enhancement principles other than the use of highly entangled states.

Many results have been obtained over the last years for such alternative schemes that are worth a comprehensive and exhaustive review that compares their usefulness with respect to the main-stream research focused on highly entangled states. We structure the review by different ways of breaking the conditions that are known to lead to SQL scaling of the sensitivity. Firstly, by going away from pure states, more general forms of quantum correlations such as quantum discord become possible. Secondly, in the derivation of the SQL the quantum systems are distinguished by an index i , which supposes that they are distinguishable. Cold atoms, on the other hand, have to be considered in general as indistinguishable particles, and the same is true for photons, which have been used for quantum enhanced measurements from the very beginning. Hence, statements about the necessity of entanglement have to be re-examined for indistinguishable particles. Thirdly, the structure of the hamiltonian is rather restrictive: a.) many hamitonians do not have a bound spectrum characterized by largest and smallest eigenvalues Λ and λ as assumed in eq.(2). Indeed, one of the most common systems used in quantum metrology, the harmonic oscillator that represents e.g. a single mode of an electro-magnetic field, has an unbound spectrum. And b.), the hamiltonian $H(\theta) = \theta \sum_{i=1}^N h_i$ does not allow for any interactions. Fourthly, unitary evolutions with a hamiltonian that depends on the parameter are not the

only way of coding a parameter in a state. In statistical mechanics, for example, there are parameters that describe the statistical ensemble, such as temperature or chemical potential for systems in thermal equilibrium, but which are not of hamiltonian origin. The same is true for non-equilibrium states. For many of these situations, the corresponding QCRs have been obtained only recently, and it often turned out that improvements beyond the SQL should be possible.

The insight that entanglement is not necessary for breaking the SQL (Tilma *et al.*, 2010) has been accepted by a growing number of researchers, but these still appear to be a minority. The situation is comparable to other aspects and fields of quantum information treatment, where previously it was thought that entanglement is necessary. For example, for a long time entanglement has been considered as necessary for non-locality, until it was realized that certain aspects of non-locality can arise without entanglement (Bennett *et al.*, 1999). Recent reviews of QEM schemes using entanglement (Degen *et al.*, 2016; Giovannetti *et al.*, 2011; Paris, 2009; Pezzé and Smerzi, 2014; Pezzè *et al.*, 2016; Tóth and Apellaniz, 2014) are available and we do not survey this vast literature here, but focus rather exclusively on QEM schemes that are not essentially based on the use of entanglement, hoping that our review will stimulate research in these directions. Before reviewing these schemes, we give the precise definition of the QCRB.

B. Fisher information

The most widely used quantity for establishing the ultimate sensitivity of measurement schemes is the quantum Fisher information (QFI) I_θ , related to the minimal variance of any asymptotically unbiased estimator function by the QCRB (1). For a possibly mixed state $\rho(\theta)$ that depends parametrically on the parameter θ , it is defined as

$$I_\theta = \text{tr}(\rho(\theta)L(\theta)^2) = \text{tr}(\partial_\theta \rho(\theta)L(\theta)), \quad (3)$$

where $L(\theta)$ is the symmetric logarithmic derivative, defined indirectly through

$$\partial_\theta \rho(\theta) = \frac{L(\theta)\rho(\theta) + \rho(\theta)L(\theta)}{2}. \quad (4)$$

In (Braunstein and Caves, 1994) it was shown that I_θ is a geometric measure on how much $\rho(\theta)$ and $\rho(\theta + d\theta)$ differ, where $d\theta$ is an infinitesimal increment of θ . The geometric measure is given by the Bures-distance,

$$ds_{\text{Bures}}^2(\rho, \sigma) \equiv 2 \left(1 - \sqrt{F(\rho, \sigma)} \right), \quad (5)$$

where the fidelity $F(\rho, \sigma)$ is defined as

$$F(\rho, \sigma) = \|\rho^{1/2}\sigma^{1/2}\|_1^2, \quad (6)$$

and $\|A\|_1 \equiv \text{tr}\sqrt{AA^\dagger}$ denotes the trace norm (Miszczyk *et al.*, 2009). With this (Braunstein and Caves, 1994),

$$I_\theta = 4ds_{\text{Bures}}^2(\rho(\theta), \rho(\theta + d\theta))/d\theta^2, \quad (7)$$

unless the rank of $\rho(\theta)$ changes with θ and thus producing removable singularities (Šafránek, 2016), a situation which do not consider in this review. The QCRB thus offers the physically intuitive picture that the parameter θ can be measured the more precisely the more strongly the state $\rho(\theta)$ depends on it. The Bures-distance is in general difficult to calculate for mixed states, but for pure states $\rho(\theta) = |\psi(\theta)\rangle\langle\psi(\theta)|$, the QFI reduces to the overlap of the derivative of the state $|\partial_\theta\psi(\theta)\rangle$ with itself and the original state, $I_\theta = 4(\langle\partial_\theta\psi(\theta)|\partial_\theta\psi(\theta)\rangle + \langle\partial_\theta\psi(\theta)|\psi(\theta)\rangle^2)$ (Paris, 2009).

In principle, the optimal measurement that saturates the QCRB can be constructed by diagonalizing $L(\theta)$. The projectors onto its eigenstates form a POVM that yields the optimal measurement. However, such a construction requires that the precise value of the parameter θ is already known. If that was the case, one could skip the measurement altogether and choose the estimator as $\theta_{\text{est}} = \theta$, with vanishing uncertainty, i.e. apparently violating the QCRB in most cases (Chapeau-Blondeau, 2015) (note, however, that for a state that depends on θ , the condition $\theta_{\text{est}} = \theta$ for an unbiased estimator cannot be fulfilled in a whole ϵ -interval about θ , such that there is no formal contradiction). If θ is not known, the more common approach is therefore to use the QCRB as a benchmark as function of θ , and then check whether physically motivated measurements can achieve it. More general schemes have been proposed to mitigate the problem of prior knowledge of the parameter. This includes the van Trees inequality (Gill and Levit, 1995; van Trees, 2001), Bayesian approaches (Macieszczak *et al.*, 2014; Rivas and Luis, 2012), adaptive measurements (Armen *et al.*, 2002; Berry and Wiseman, 2000, 2002, 2006, 2013; Fujiwara, 2006; Higgins *et al.*, 2009; Okamoto *et al.*, 2012; Serafini, 2012; Wheatley *et al.*, 2010; Wiseman, 1995), and approaches specialized to particular parameter estimation problems such as phase estimation (Hall *et al.*, 2012). Another point to be kept in mind is that the QCRB can be reached asymptotically for a large number of measurements, but not necessarily for a finite number of measurements. The latter case is clearly relevant for experiments and subject of active current research (see e.g. (Liu and Yuan, 2016)).

These limitations notwithstanding, we base this review almost exclusively on the QCRB (with the exception of Sec.II.F on quantum channel discrimination and parts of Sec.IV.C.1, where a signal-to-noise ratio is used), given that the overwhelming majority of results have been obtained for it and allow an in-depth comparison of different strategies.

Two more properties of the QFI are worth mentioning: $I_\theta(\rho(\theta))$ is a convex function of $\rho(\theta)$, i.e. for two density matrices $\rho(\theta)$ and $\sigma(\theta)$ and $0 \leq \lambda \leq 1$ we have (Fujiwara, 2001)

$$I_\theta(\lambda\rho(\theta) + (1-\lambda)\sigma(\theta)) \leq \lambda I_\theta(\rho(\theta)) + (1-\lambda)I_\theta(\sigma(\theta)). \quad (8)$$

And the QFI of uncorrelated states is additive, (Fujiwara and Hashizumé, 2002):

$$I_\theta(\rho(\theta) \otimes \sigma(\theta)) = I_\theta(\rho(\theta)) + I_\theta(\sigma(\theta)). \quad (9)$$

Convexity implies that the precision of measurements cannot be increased by classically mixing states with mixing probabilities independent of the parameter (Braun, 2010). Additivity immediately leads to $1/\sqrt{N}$ scaling in eq.(2), as the QFI of the N uncorrelated subsystems is just N times the QFI of a single subsystem, and hence no measurement of the whole system, even if entangling, can improve the sensitivity when the parameter was already imprinted on a product state.

The Fisher information can be generalized to multi-parameter estimation (Helstrom, 1969; Paris, 2009), $\boldsymbol{\theta} = (\theta_1, \theta_2, \dots)$. The Bures distance between two infinitesimal close states then reads

$$ds_{\text{Bures}}^2(\rho_{\boldsymbol{\theta}}, \rho_{\boldsymbol{\theta}+d\boldsymbol{\theta}}) = 2 \left(1 - \text{tr} \sqrt{\sqrt{\rho_{\boldsymbol{\theta}}} \rho_{\boldsymbol{\theta}+d\boldsymbol{\theta}} \sqrt{\rho_{\boldsymbol{\theta}}}} \right). \quad (10)$$

An expansion of $ds_{\text{Bures}}^2(\rho_{\boldsymbol{\theta}}, \rho_{\boldsymbol{\theta}+d\boldsymbol{\theta}})$ leads to the quantum Fisher information matrix (Paris, 2009; Sommers and Zyczkowski, 2003),

$$ds_{\text{Bures}}^2(\rho_{\boldsymbol{\theta}}, \rho_{\boldsymbol{\theta}+d\boldsymbol{\theta}}) = \frac{I_{\theta_k, \theta_{k'}}}{4} d\theta_k d\theta_{k'}. \quad (11)$$

The QCRB generalizes to a lower bound on the covariance matrix $\text{Cov}[\boldsymbol{\theta}]$ of the parameters θ_i (Helstrom, 1969, 1976; Paris, 2009),

$$\text{Cov}[\boldsymbol{\theta}] \geq \frac{1}{M} (I(\boldsymbol{\theta}))^{-1}, \quad (12)$$

where $\text{Cov}[\boldsymbol{\theta}]_{ij} = \langle\theta_i\theta_j\rangle - \langle\theta_i\rangle\langle\theta_j\rangle$, and $A \geq B$ means that $A - B$ is a positive-semidefinite matrix. Contrary to the single parameter QCRB, the bound (12) can in general not be saturated, even in the limit of infinitely many measurements. The Bures metric has also been called fidelity susceptibility in the framework of quantum phase transitions (QPTs) (Gu, 2010).

II. QUANTUM CORRELATIONS BEYOND ENTANGLEMENT

A. Parallel versus sequential strategies in unitary quantum metrology

One of the most typical applications of quantum metrology is the task of unitary parameter estimation, exemplified in particular by phase estimation

(Giovannetti *et al.*, 2006, 2011). Let $U_\theta = \exp(-i\theta H)$ be a unitary transformation, with θ the unknown parameter to be estimated, and H a selfadjoint Hamiltonian operator which represents the generator of the transformation. The typical estimation procedure then consists of the following steps: a) preparing an input probe in a state ρ ; b) propagating the state with the unitary transformation U_θ ; c) measuring the output state $\rho_\theta = U_\theta \rho U_\theta^\dagger$; d) performing classical data analysis to infer an estimator θ_{est} for the parameter θ .

Let us now assume one has the availability of N utilizations of the transformation U_θ . Then, the use of N uncorrelated probes in a global initial state $\rho^{\otimes N}$, each of which is undergoing the transformation U_θ in parallel, yields an estimator whose minimum variance scales as $1/N$ (SQL). On the other hand, by using an initial entangled state ρ of the N probes, and propagating each with the unitary U_θ in parallel, one can in principle achieve the HL, meaning that an optimal estimator θ_{est} can be constructed whose asymptotic variance, in the limit $N \gg 1$, scales as $1/N^2$. However, it is not difficult to realize that the very same precision can be reached without the use of entanglement, by simply preparing a single input probe in a superposition state with respect to the eigenbasis of the generator H , and letting the probe undergo N sequential iterations of the transformation U_θ .

For instance, thinking of each probe as a qubit for simplicity, and fixing the generator H to be the Pauli matrix σ_z , one can either consider a parallel scheme with N input probes in the GHZ (or cat-like) maximally entangled state $|\Psi\rangle = (|00\dots 0\rangle + |11\dots 1\rangle)/\sqrt{2}$, or a sequential scheme with a single probe in the superposition $|\psi\rangle = (|0\rangle + |1\rangle)/\sqrt{2}$: both schemes reach the HL scaling in the estimation of the phase shift θ , meaning that the quantum Cramér-Rao bound can be asymptotically saturated in both cases by means of an optimal measurement, associated to a QFI scaling quadratically with N . The equivalence between entanglement in parallel schemes and *coherence* (namely, superposition in the eigenbasis of the generator) (Baumgratz *et al.*, 2014; Streltsov *et al.*, 2016) in sequential schemes further extends to certain schemes of quantum metrology in the presence of noise, namely when the unitary encoding the parameter to be estimated and the noisy channel commute with each other (e.g. in the case of phase estimation affected by dephasing) (Boixo and Heunen, 2012; Demkowicz-Dobrzański and Maccone, 2014), although in more general instances entanglement is shown to provide an advantage (Demkowicz-Dobrzański and Maccone, 2014; Escher *et al.*, 2011; Huelga *et al.*, 1997).

Similarly, in continuous variable optical interferometry (Caves, 1981), equivalent performances can be reached (for unitary phase estimation) by using either a two-mode entangled probe, such as a N00N state, or a single-mode non-classical state, such as a squeezed state. These are elementary examples of QEM achievable without

entanglement, yet exploiting genuinely quantum effects such as nonclassicality and superposition. Such features can be understood by observing that both optical nonclassicality in infinite-dimensional systems and coherence in finite-dimensional systems can be converted to entanglement within a well-defined resource-theoretic framework (Asbóth *et al.*, 2005; Killoran *et al.*, 2016; Streltsov *et al.*, 2015; Vogel and Sperling, 2014), and can be thought-of as equivalent resources to entanglement for certain practical purposes, as is evidently the case for unitary metrology.

B. General results on the usefulness of entanglement

More generally, for unitary metrology with multipartite probes in a parallel setting, a quite general formalism has been developed to identify the metrologically useful correlations in the probes in order to achieve QEM (Pezzè and Smerzi, 2009) (see (Pezzè and Smerzi, 2014; Tóth and Apellaniz, 2014) for recent reviews). Specifically, let us consider an input state ρ of N qubits and a linear interferometer with Hamiltonian generator given by $H = J_l = \frac{1}{2} \sum_{i=1}^N \sigma_l^{(i)}$, i.e. a component of the collective (pseudo-)angular momentum of the N probes in the direction $l = x, y, z$, with $\sigma_l^{(i)}$ denoting the l th Pauli matrix for qubit i . If ρ is k -producible, i.e., it is a convex mixture of pure states which are tensor products of at most k -qubit states, then the QFI is bounded above as follows (Hyllus *et al.*, 2012; Tóth, 2012),

$$I_\theta(\rho, J_l) \leq nk^2 + (N - nk)^2, \quad (13)$$

where n is the integer part of N/k . This means that genuine multipartite entangled probes ($k = N$) are required to reach the maximum sensitivity, given by the HL $I_\theta \propto N^2$, even though partially entangled states can still result in QEM beyond the SQL.

A similar conclusion has been reached in (Augusiak *et al.*, 2016) considering the geometric measure of entanglement, which quantifies how far ρ is from the set of fully separable (1-producible) states according to the fidelity. Namely, for unitary metrology with N parallel probes initialized in the mixed state ρ , in the limit $N \rightarrow \infty$ a nonvanishing value of the geometric measure of entanglement of ρ is necessary for the exact achievement of the HL. However, a sensitivity arbitrarily close to the HL, $I_\theta \propto N^{2-\epsilon}$ for any $\epsilon > 0$, can still be attained even if the geometric measure of entanglement of ρ vanishes asymptotically for $N \rightarrow \infty$. In deriving these results, the authors proved an important continuity relation for the QFI in unitary dynamics (Augusiak *et al.*, 2016).

C. Role of quantum discord in parameter estimation with mixed probes

Here we will focus our attention on possible advantages stemming from the use of quantum correlations more general than entanglement in the (generally mixed) state of the input probes for a metrological task. Such correlations are usually referred to under the collective name of *quantum discord* (Henderson and Vedral, 2001; Modi *et al.*, 2012; Ollivier and Zurek, 2001). A state ρ of a bipartite system AB has nonzero discord (from the point of view of subsystem A) iff, by definition, it is altered by all possible local measurements performed on subsystem A . Every entangled state is also discordant, but the converse is not true; in fact, almost all separable states still exhibit nonzero discord (Ferraro *et al.*, 2010). The only bipartite states with zero discord, from the point of view of subsystem A , are so-called classical-quantum states, which take the form

$$\chi^{(A)} = \sum_i p_i |i\rangle\langle i|^A \otimes \tau_i^B, \quad (14)$$

where the states $\{|i\rangle^A\}$ form an orthonormal basis for subsystem A , and $\{\tau_i^B\}$ denote a set of arbitrary states for subsystem B , while $\{p_i\}$ stands for a probability distribution. For any bipartite state ρ , the original measure of quantum discord is defined as (Henderson and Vedral, 2001; Ollivier and Zurek, 2001)

$$D^{(A)}(\rho) = \inf_{\Pi^A} \{I(\rho) - I(\Pi^A[\rho])\}, \quad (15)$$

where $\Pi^A[\rho]$ denotes the conditional state obtained after a POVM measurement on subsystem A , and $I(\rho) = S(\text{Tr}_A \rho) + S(\text{Tr}_B \rho) - S(\rho)$ stands for the mutual information, with $S(\rho) = -\text{Tr}(\rho \log \rho)$ the von Neumann entropy; see (Adesso *et al.*, 2016; Modi *et al.*, 2012) for more details. $D^{(A)}(\rho)$ vanishes if and only if ρ is a classical-quantum state of the form (14), and in general quantifies the minimum loss in total correlations of the state ρ due to a least disturbing local measurement on A .

In a multipartite setting, one can define fully classical states as the states with zero discord with respect to all possible subsystems, or alternatively as the states which are left invariant by a set of local measurements performed on all subsystems. Such states take the form $\chi = \sum_{i_1, \dots, i_N} p_{i_1, \dots, i_N} |i_1\rangle\langle i_1|^{A_1} \otimes \dots \otimes |i_N\rangle\langle i_N|^{A_N}$ for an N -particle system $A_1 \dots A_N$; i.e., they are diagonal in a local product basis. One can think of these states as the only ones which are completely classically correlated, that is, completely described by a classical multivariate probability distribution $\{p_{i_1, \dots, i_N}\}$, embedded into a density matrix formalism. An alternative way to quantify discord in a (generally multipartite) state ρ is then by taking the distance between ρ and the set of classically correlated states, according to a suitable (quasi)distance

function. For instance, the relative entropy of discord (Modi *et al.*, 2010) is defined as

$$D_R(\rho) = \inf_{\chi} S(\rho \| \chi), \quad (16)$$

where the minimization is over all classically correlated states χ , and $S(\rho \| \chi) = \text{Tr}(\rho \log \rho - \rho \log \chi)$ denotes the quantum relative entropy. For a dedicated review on measures of discord-type quantum correlations see (Adesso *et al.*, 2016).

In a metrological context, (Modi *et al.*, 2011) investigated unitary phase estimation using N -qubit probe states initialized in mixed states with a) no correlations; b) only classical correlations; or c) quantum correlations (discord and/or entanglement). All the families of probe states had the same spectrum, i.e. in particular the same degree of mixedness (which is a meaningful assumption if one is performing a metrology experiment in an environment with a fixed common temperature), and were chosen due to their relevance in recent NMR experiments. In particular, given an initial thermal state $\rho_0(p) = (\frac{1+p}{2}|0\rangle\langle 0| + \frac{1-p}{2}|1\rangle\langle 1|)$ for each single qubit (with purity parameter $0 \leq p \leq 1$), the product states $\rho_N^{(a)}(p) = [H\rho_0(p)H]^{\otimes N}$ were considered for case a), and the GHZ-diagonal states $\rho_N^{(c)}(p) = CH_1C\rho_0(p)^{\otimes N}CH_1C$ were considered for case c), with H denoting the single-qubit Hadamard gate (acting on each qubit in the first case, and only on the first qubit in the second case), and $C = \otimes_{j=2}^N \text{C-NOT}_{1j}$ a series of Control-Not operations acting on pairs of qubits 1 and j . The two considered classes of states result in QFIs $I_\theta^{(a)} = p^2N$ and $I_\theta^{(c)} \gtrsim p^2N^2$, respectively. Therefore the authors of (Modi *et al.*, 2011) concluded that QEM with an enhancement scaling factor $I_\theta^{(c)}/I_\theta^{(a)} \approx N$ is possible using pairs of mixed probe states with arbitrary degree of purity. This enhancement persists even when the states in strategy c) are fully separable, which occurs for $p \lesssim a + b/N$ (with a, b determined numerically for each N), in which case both strategies perform sub-SQL, yet the quadratic enhancement of c) over a) is maintained, as it is independent of p . The authors then argue that multipartite quantum discord — which increases with N according to the relative entropy measure of Eq. (16) and vanishes only at $p = 0$ — may be responsible for this enhancement. Let us remark that, even though the QFI is convex (which means that for every separable but discordant mixed state there exists a pure product state with a higher or equal QFI), the analysis in (Modi *et al.*, 2011) was performed at fixed spectrum (and thus degree of purity) of the input probes, a constraint which allowed the authors to still identify an advantage in using correlations weaker than entanglement, as opposed to no correlations. However, it is presently unclear whether these conclusions are special to the selected classes of states, or can be further extended to more general settings.

In a more recent work, (Cable *et al.*, 2016) analyzed a model of unitary quantum metrology inspired by the computational algorithm known as deterministic quantum computation with one quantum bit (DQC1) or “power of one-qubit” (Knill and Laflamme, 1998). Using only one pure qubit supplemented by a register of l maximally mixed qubits, all individually subject to a local unitary phase shift U_θ , their model was shown to achieve the SQL for the estimation of θ , which can be conventionally obtained using the same number of qubits in pure uncorrelated states. They found that the SQL can be exceeded by using one additional qubit, which only contributes a small degree of extra purity, which, however, for any finite amount of extra purity leads to an entangled state at the stage of parameter encoding. In this model, incidentally, the output state after the unitary encoding was found to be always separable but discordant, with its discord vanishing only in the limit of vanishing variance of the estimator for the parameter θ . It is not quite clear if and how the discord in the final state can be interpreted in terms of a resource for metrology, but the achievement of the SQL with all but one probes in a fully mixed state was identified as a quantum enhancement without the use of entanglement. This suggests that further investigation on the role of quantum discord (as well as state purity) in metrological algorithms with vanishing entanglement may be in order. A protocol for multiparameter estimation using DQC1 was studied in (Boixo and Somma, 2008), although the resource role of correlations was not discussed there. In (MacCormick *et al.*, 2016) a detailed investigation of a DQC1-based protocol was made based on coherently controlled Rydberg interactions between a single atom and an atomic ensemble containing N atoms. The protocol allows one to estimate a phase shift assumed identical for all atoms in the atomic ensemble with a sensitivity that interpolates smoothly between SQL and HL when the purity of the atomic ensemble increases from a fully mixed state to pure states. It leads to a cumulative phase shift proportional to N , and the scheme can in fact also be seen as an implementation of “coherent averaging”, with the control qubit playing the role of the “quantum bus” (see Sec. IV.H).

D. Black-box metrology and the interferometric power

As explicitly discussed in Sec. II.A, for unitary parameter estimation, if one has full prior information on the generator H of the unitary transformation U_θ encoding the parameter θ , then no correlations are required whatsoever, and probe states with coherence in the eigenbasis of H suffice to achieve QEM in a sequential scheme. Recently, (Adesso, 2014; Girolami *et al.*, 2014, 2013) investigated quantum metrology in a black-box paradigm, where the generator H is assumed not fully known a priori. In such a case, suppose one fixes an input single-

particle probe ρ , then it is impossible to guarantee a precision in the estimation of θ for all possible nontrivial choices of H . This is because, in the worst case scenario, the black-box unitary transformation may be generated by a H which commutes with the input state ρ , resulting in no information imprinted on the probe, hence in a vanishing QFI. It is clear then that, to be able to estimate parameters independently of the generator, one needs an ancillary system correlated with the probe. It is in this precise context that discord-type correlations, rather than entanglement or classical correlations, are found to play the resource role.

Consider a simple two-arm interferometric configuration, and let us retrace the steps of parameter estimation in the black-box setting (Girolami *et al.*, 2014): a) knowing only the spectrum of H , but not its eigenbasis, N copies of an input state ρ of two particles, the probe A and the ancilla B , are prepared; b) in each iteration, particle A enters a black-box where it undergoes a unitary transformation $U_\theta = \exp(-i\theta H)$, while particle B is transmitted with no interaction; c) the agent controlling the black-box announces the generator H , so that parties A and B can jointly perform the best possible measurement on the N copies of the two-particle output state $\rho_\theta = (U_\theta \otimes \mathbb{I})\rho(U_\theta \otimes \mathbb{I})^\dagger$; d) an optimal unbiased estimator θ_{est} is constructed for the parameter θ . For any given black-box setting H , the QFI $I_\theta(\rho, H)$ determines the maximal precision enabled by the input state ρ in estimating the parameter θ generated by H , via the asymptotic saturation of the quantum Cramér-Rao bound.

One can then introduce a figure of merit quantifying the worst case precision, guaranteed by the state ρ . This is done naturally by minimizing the QFI over all generators H with a given spectral class (the spectrum is assumed nondegenerate, with a canonical choice being that of equispaced eigenvalues) (Girolami *et al.*, 2014, 2013). The *interferometric power* (IP) of the bipartite state ρ with respect to the probing system A is then defined as

$$P^{(A)}(\rho) = \frac{1}{4} \min_H I_\theta(\rho, H). \quad (17)$$

As proven in (Girolami *et al.*, 2014), the IP is a measure of discord-type correlations in the input state ρ . In particular, it vanishes if and only if ρ takes the form of a classical-quantum state, Eq. (14). This entails that states with zero discord cannot guarantee a precision in parameter estimation in the worst case scenario, while any other bipartite state (entangled or separable) with finite discord is suitable for estimating parameters encoded by a unitary transformation (acting on one subsystem) no matter the generator, with minimum guaranteed precision quantified by the IP of the state. This conclusion holds both for parameter estimation in finite-dimensional systems (Girolami *et al.*, 2014), and for continuous-variable optical interferometry (Adesso, 2014). Recently, it has been shown more for-

mally that entanglement accounts only for a portion of the quantum correlations relevant for bipartite quantum interferometry. In particular, the IP, which is by definition a lower bound to the QFI (for any fixed generator H), is itself bounded from below (for bipartite systems of any dimension) by a measure of entanglement named the interferometric entanglement, which reduces to the squared concurrence for two-qubit states (Bromley *et al.*, 2016). The IP can be evaluated in closed form, solving analytically the minimization in Eq. (17), for all finite-dimensional states such that subsystem A is a qubit (Girolami *et al.*, 2014), and for all two-mode Gaussian states when the minimization is restricted to Gaussian unitaries (Adesso, 2014). An experimental demonstration of black-box QEM relying on discordant states as opposed to classically correlated states has been reported using a two-qubit NMR ensemble realized in chloroform (Girolami *et al.*, 2014).

E. Estimation of loss in Gaussian quantum metrology

Any quantum optical communication, from fibre-based to free-space implementations, is inevitably affected by energy dissipation. The fundamental model to describe this scenario is the lossy channel. This attenuates an incoming bosonic mode by transmitting a fraction $\eta \leq 1$ of the input photons, while sending the other fraction $1 - \eta$ into the environment. The maximum number of bits per channel use at which we can transmit quantum information, distribute entanglement or generate secret keys through such a lossy channel are all equal to $-\log(1 - \eta)$ (Pirandola *et al.*, 2015), a fundamental rate-loss tradeoff that only quantum repeaters may surpass (Pirandola, 2016). For these and other implications to quantum communication, it is of paramount importance to estimate the transmissivity of a lossy channel in the best possible way.

Quantum estimation of bosonic loss was first studied in Ref. (Monras and Paris, 2007) by using single-mode pure Gaussian states (see also Ref. (Pinel *et al.*, 2013)). In this setting, the performance of the coherent state probes at fixed input energy provides the shot-noise limit or classical benchmark, which has to be beaten by truly quantum probes. Let us denote by \bar{n} the mean number of photons, then the shot-noise limit is equal to $I_\eta \simeq \eta^{-1}\bar{n}$ (Monras and Paris, 2007; Pinel *et al.*, 2013). The use of squeezing can beat this limit, following the original intuition for phase estimation of Ref. (Caves, 1981). In fact, Ref. (Monras and Paris, 2007) showed that, in the regime of small loss $\eta \simeq 1$ and small energy $\bar{n} \simeq 0$, the optimal input state is a squeezed vacuum state and this reaches the optimal QFI as a function of \bar{n} , which is $I_\eta \simeq 4\bar{n}$. The use of squeezing for estimating the interaction parameter in bilinear bosonic Hamiltonians (including beam-splitter interactions) was also dis-

cussed in Ref. (Gaiba and Paris, 2009), showing that unentangled single-mode squeezed probes offer equivalent performance to entangled two-mode squeezed probes for practical purposes.

The optimal scaling $I_\eta \simeq 4\bar{n}$ is also achieved by using Fock states at the input (Adesso *et al.*, 2009). This is the best unconditional strategy, in the sense that its performance does not depend on the value of the transmissivity η . It also beats adaptive strategies based on coherent states, whenever $\eta \geq 1/4$. Note that, because Fock states can only be used when the input energy corresponds to integer photon numbers, in all the other cases one needs to engineer superpositions, e.g., between the vacuum and the one-photon Fock state if we want to explore the low-energy regime $\bar{n} \lesssim 1$. Non-Gaussian qutrit and quartet states can be designed to beat the best Gaussian probes (Adesso *et al.*, 2009). It is still an open question to determine the optimal probes for estimating loss at any energy regime. We certainly know the bound $I_\eta \leq 4\bar{n}$ for any \bar{n} , which can be proven by dilating the lossy channel into a beamsplitter unitary and then performing parameter estimation (Monras and Paris, 2007). It is also an open question to find the best input states in the presence of adaptive strategies.

Interestingly, the problem of estimating the loss parameters of a pair of lossy bosonic channels has been proven formally equivalent to the problem estimating the separation of two incoherent optical point-like sources (Lupo and Pirandola, 2016). In this context (Tsang *et al.*, 2016) showed that a pair of highly-attenuated thermal sources can be resolved independently from their separation if one adopts quantum measurements based on photon counting, instead of standard intensity measurements. Thus, quantum detection strategies enables one to beat the so-called ‘‘Rayleigh’s curse’’ which affects classical imaging (Tsang *et al.*, 2016). This curse is reinstated in the classical limit of very bright thermal sources (Lupo and Pirandola, 2016; Nair, R. and Tsang, T., 2016). On the other hand, (Lupo and Pirandola, 2016) showed that quantum-correlated sources can be super-resolved at the sub-Rayleigh scale. In fact, it is possible to engineer quantum-correlated point-like sources that are not entangled (but discordant) which displays super-resolution, so that the closer the sources are the better their distance can be estimated.

The estimation of loss becomes complicated in the presence of decoherence, such as thermal noise in the environment and non-unit efficiency of the detectors. From this point of view, Ref. (Spedalieri *et al.*, 2016) considered a very general model of Gaussian decoherence which also includes the potential presence of non-Markovian memory effects. In such a scenario, Ref. (Spedalieri *et al.*, 2016) showed the utility of asymmetrically correlated thermal states (i.e., with largely different photon numbers in the two modes), fully based on

discord and void of entanglement. These states can be used to estimate bosonic loss with a sensitivity that approaches the shot noise limit and may also surpass it in the presence of correlated noise and memory effects in the environment. This kind of thermal quantum metrology has potential applications for practical optical instruments (e.g., photometers) or at different wavelengths (e.g., far infrared, microwave or X-ray) where the generation of quantum features, such as coherence, number states, squeezing or entanglement, may be challenging.

Clearly we may also consider the estimation of other parameters in bosonic Gaussian channels, such as thermal or additive noise (Pirandola, S. and Lupo, C., 2016). Loss and thermal noise may also be inferred jointly (Bellomo *et al.*, 2009, 2010a,b; Monras and Illuminati, 2011). In general, Gaussian quantum metrology aims at estimating any parameter encoded in a Gaussian channel. If we employ Gaussian states at the input, then we have as output Gaussian states and we may exploit closed formulas for the QFI. These formulas can be derived by direct evaluation of the symmetric logarithmic derivative (SLD) (Monras, 2013) or by considering the infinitesimal expression of the quantum fidelity (Banchi *et al.*, 2015; Pinel *et al.*, 2012, 2013). The latter approach may exploit the general formula of Ref. (Banchi *et al.*, 2015). In fact, for two arbitrary multimode Gaussian states, ρ_1 and ρ_2 , with mean values u_1 and u_2 , and covariance matrices V_1 and V_2 , we may write the Uhlmann-Jozsa fidelity (Banchi *et al.*, 2015)

$$F(\rho_1, \rho_2) = \frac{F_{\text{tot}} e^{-\frac{1}{2}(u_2 - u_1)^T (V_1 + V_2)^{-1} (u_2 - u_1)}}{\sqrt{\det(V_1 + V_2)}}, \quad (18)$$

$$F_{\text{tot}}^2 = \det \left[2 \left(\sqrt{\mathbb{1} + \frac{(V_{\text{aux}} \Omega)^{-2}}{4}} + \mathbb{1} \right) V_{\text{aux}} \right] \quad (19)$$

where we set $V_{\text{aux}} := \Omega^T (V_1 + V_2)^{-1} (\Omega/4 + V_2 \Omega V_1)$ with Ω being the symplectic form (Banchi *et al.*, 2015). Specific expressions for the fidelity were previously given for single-mode Gaussian states (Scutaru, 1998), two-mode Gaussian states (Marian and Marian, 2012), and multimode Gaussian states assuming that one of the states is pure (Spedalieri *et al.*, 2013).

From Eq. (18) we may derive the Bures metric ds^2 . In fact, consider two infinitesimally-close Gaussian states ρ , with statistical moments u and V , and $\rho + d\rho$, with statistical moments $u + du$ and $V + dV$. Expanding at the second order in du and dV , one finds (Banchi *et al.*, 2015)

$$ds^2 = 2[1 - \sqrt{F(\rho, \rho + d\rho)}] = \frac{du^T V^{-1} du}{4} + \frac{\delta}{8}, \quad (20)$$

where $\delta := 4\text{Tr}[dV(4\mathcal{L}_V + \mathcal{L}_\Omega)^{-1}dV]$, $\mathcal{L}_A X := AXA$, and the inverse of the superoperator $4\mathcal{L}_V + \mathcal{L}_\Omega$ refers to the pseudo-inverse. A similar expression was also computed

by (Monras, 2013) using the SLD. From the Bures metric we may directly derive the QFI (see eq.(7)), and simple analytical formulas are known for pure multi-mode Gaussian states.

In this context, Ref. (Pinel *et al.*, 2012) determined the QCRB for estimating a parameter which is directly encoded in a pure multi-mode Gaussian state. It was realized that, in the limit of large photon number, no entanglement nor correlations between different modes are necessary for obtaining the optimal sensitivity. Rather, a detection mode can be used based on the derivative of the mean photon field with respect to the parameter θ , into which all the resources in terms of photons and squeezing should be put. The mean photon field is defined as $\bar{a}_\theta(\mathbf{r}, t) = \langle \psi_\theta | a(\mathbf{r}, t) | \psi_\theta \rangle$, with all parameter dependence in the pure Gaussian quantum state $|\psi\rangle_\theta$, $a(\mathbf{r}, t) = \sum_i a_i v_i(\mathbf{r}, t)$, where $v_i(\mathbf{r}, t)$ are orthonormal mode functions found from solving Maxwell's equation with appropriate boundary conditions, a_i is the annihilation operators of mode i , and the sum is over all modes. The mean field can be normalized, $u_\theta = \bar{a}_\theta(\mathbf{r}, t)/\|\bar{a}_\theta\|$, where the norm $\|f(\mathbf{r}, t)\| = (\int |f(\mathbf{r}, t)|^2 d^2\mathbf{r} dt)^{1/2}$ contains spatial integration over a surface perpendicular to the light beam propagation and temporal integration over the detection time. The detection mode is then defined as $\tilde{v}_1(\mathbf{r}, t) = \frac{\bar{a}'_\theta(\mathbf{r}, t)}{\|\bar{a}'_\theta\|}$ and leads to a QFI

$$I_\theta = N_\theta \left(4\|u'_\theta\|^2 + \left(\frac{N'_\theta}{N_\theta} \right)^2 \right) V_{\theta, [1,1]}^{-1}, \quad (21)$$

where N_θ is the mean photon number, and $V_{\theta, [1,1]}^{-1}$ the matrix element of the inverse covariance matrix of the Gaussian state corresponding to the detection mode. All other modes are chosen orthonormal to it. The SQL corresponds to a QFI of a coherent state, in which case $V_{\theta, [1,1]}^{-1} = 1$. Hence, an improvement over the SQL is possible with pure Gaussian states by squeezing the detection mode. The scaling with N_θ can be modified if $V_{\theta, [1,1]}^{-1}$ depends on N_θ . For a fixed total energy a scaling $I_\theta \propto N_\theta^{3/2}$ can be achieved. This was proposed in (Barnett *et al.*, 2003) for measuring a beam displacement. The QCRB in eq.(21) can be reached by homodyne detection with the local oscillator in this detection mode.

F. Quantum channel discrimination

A fundamental protocol which is closely related to quantum metrology is quantum channel discrimination (Acin, 2001; Childs *et al.*, 2000; Invernizzi *et al.*, 2011; Lloyd, 2008; Pirandola, 2011; Sacchi, 2005; Tan *et al.*, 2008), which may be seen as a sort of digitalized version of quantum metrology. Its basic formulation is binary and involves the task of distinguishing between two quantum channels, \mathcal{E}_0 or \mathcal{E}_1 , associated with two a

priori probabilities $\pi_0 := \pi$ and $\pi_1 = 1 - \pi$. During the encoding phase, one of such channels is picked by Alice and stored in a box which is then passed to Bob. In the decoding phase, Bob uses a suitable state at the input of the box and performs a quantum measurement of its output. Bob may also use ancillary systems which are quantum correlated with the input probes and are directly sent to the measurement. For the specific tasks of discriminating bosonic channels, the input is assumed to be constrained in energy, so that we fix the mean number of photons \bar{n} per input probe, or more strongly, the mean total number of photons which are globally irradiated through the box (Weedbrook *et al.*, 2012).

Quantum channel discrimination is an open problem in general. However, when we fix the input state, it is translated into an easier problem to solve, i.e., the quantum discrimination of the output states. In the binary case, this conditional problem has been fully solved by the so-called Helstrom bound which provides the minimum mean error probability \bar{p} in the discrimination of any two states ρ_0 and ρ_1 . Assuming equiprobable states ($\pi = 1/2$), this bound is simply given by their trace distance D , i.e., we have (Helstrom, 1976)

$$\bar{p} = \frac{1}{2} [1 - D(\rho_0, \rho_1)]. \quad (22)$$

In the case of multi-copy discrimination, where we probe the box N times and we aim to distinguish the two outputs $\rho_0^{\otimes N}$ and $\rho_1^{\otimes N}$, the mean error probability $\bar{p}(N)$ may be not so easy to compute and, therefore, we resort to suitable bounds. Using the quantum fidelity $F(\rho_0, \rho_1)$ from (6), and setting

$$Q(\rho_0, \rho_1) := \inf_{0 \leq s \leq 1} \text{Tr}(\rho_0^s \rho_1^{1-s}), \quad (23)$$

we may then write (Audenaert *et al.*, 2007; Fuchs and de Graaf, 1999)

$$\frac{1 - \sqrt{1 - F(\rho_0, \rho_1)^N}}{2} \leq \bar{p}(N) \leq \frac{Q^N(\rho_0, \rho_1)}{2}, \quad (24)$$

where $Q^N(\rho_0, \rho_1)/2$ is the quantum Chernoff bound (QCB) (Audenaert *et al.*, 2007). In particular, the QCB is asymptotically tight for large N . Furthermore, it can be easily computed for arbitrary multimode Gaussian states (Pirandola and Lloyd, 2008).

Since the conditional output states can be optimally distinguished, the non-trivial part in quantum channel discrimination is the optimization of the mean error probability \bar{p} over the input states. For this reason, it is an extremely rich problem and depending on the types of quantum channels, quantum correlations may play an important role or not. We now discuss some specific cases in more detail.

Quantum channel discrimination has various practical applications. One which is very well known is quantum

illumination (Lloyd, 2008; Tan *et al.*, 2008) which forms the basis for a “quantum radar” (Barzanjeh *et al.*, 2015). Despite the fact that entanglement is used at the input between the signal (sent to probe a potential target) and the idler (kept at the radar state for joint detection), entanglement is completely absent at the output between reflected and idler photons. Nevertheless the scheme assures a superior performance with respect to the use of coherent states; in particular, an increase by a factor 4 of the exponent $-(\ln P(e))/M$ of the asymptotic error probability $P(e)$ (where M is the number of transmissions) (Tan *et al.*, 2008). For this reason, the quantum illumination advantage has been studied in relation with the consumption of other discord-type quantum correlations beyond entanglement (Bradshaw *et al.*, 2016; Weedbrook *et al.*, 2016). More precisely, the enhanced performance of quantum illumination (with respect to signal probing not assisted by an idler) corresponds to the amount of discord which is expended to resolve the target (i.e., to encode the information about its presence or absence).

Another application of quantum channel discrimination is quantum reading (Pirandola, 2011). Here the basic aim is to discriminate between two different channels which are used to encode an information bit in a cell of a classical memory. In an optical setting, this means to discriminate between two different reflectivities, generally assuming the presence of decoherence effects, such as background stray photons. The maximum amount of bits per cell that can be read is called “quantum reading capacity” (Pirandola *et al.*, 2011). This model has also been studied in the presence of thermal and correlated decoherence, as that arising from optical diffraction (Lupo *et al.*, 2013). In all cases, the classical benchmark associated with coherent states can be largely beaten by non-classical states, as long as the mean number of photons hitting the memory cells is suitably low.

Depending on the regime, we may choose a different type of non-classical states. In the presence of thermal decoherence induced by background photon scattering, two-mode squeezed vacuum states between signal modes (reading the cells) and idler modes (kept for detection) are nearly-optimal. However, in the absence of decoherence, the sequential readout of an ideal memory (where one of the reflectivities is exactly 100%) is optimized by number states at the input (Nair, 2011). Ref. (Roga *et al.*, 2015) showed that, in specific regimes, the quantum advantage can be related with a particular type of quantum correlations, the discord of response, which is defined as the trace, or Hellinger, or Bures minimum distance from the set of unitarily perturbed states (Roga *et al.*, 2014). Ref. (Roga *et al.*, 2015) also identified particular regimes where strongly discordant states are able to outperform pure entangled transmitters.

Let us consider the specific case of unitary channel dis-

crimination. Suppose that the task is to decide whether a unitary U_θ was applied or not to a probing subsystem A of a joint system (A, B) . In other words, the aim is to discriminate between the two possible output states $\rho_\theta = (U_\theta \otimes \mathbb{I})\rho(U_\theta \otimes \mathbb{I})^\dagger$ (when the unitary U_θ has acted on A) or ρ (equal to the input, when the identity has acted on A instead). In the limit of an asymptotically large number $N \gg 1$ of copies of ρ , the minimal probability of error in distinguishing between ρ_θ and ρ , using an optimal discrimination strategy scales approximately as the QCB $Q(\rho, \rho_\theta)^N/2$.

It is clear that the quantity $1 - Q(\rho, \rho_\theta)$ plays a similar role in the present discrimination context as the QFI in the parameter estimation scheme. One can therefore introduce an analogous figure of merit quantifying the worst case ability to discriminate, guaranteed by the state ρ . The *discriminating strength* (DS) of the bipartite state ρ with respect to the probing system A is then defined as (Farace *et al.*, 2014)

$$D^{(A)}(\rho) = \min_H [1 - Q(\rho, \rho_\theta)], \quad (25)$$

where the minimization is performed once more over all generators H within a given nondegenerate spectrum.

As proven in (Farace *et al.*, 2014), the DS is another measure of discord-type correlations in the input state ρ , which vanishes if and only if ρ is a classical-quantum state as in Eq. (14). The DS is also computable in closed form for all finite-dimensional states such that subsystem A is a qubit. In the latter case, the DS turns out to be proportional to the local quantum uncertainty (Girolami *et al.*, 2013), a further measure of discord-type correlations defined as in Eq. (17), but with the QFI replaced by the Wigner-Yanase skew information (Girolami *et al.*, 2013). The DS has also been extended to continuous-variable systems, and evaluated for special families of two-mode Gaussian states restricting the minimization in Eq. (25) to Gaussianity-preserving generators (i.e., quadratic Hamiltonians) (Rigovacca *et al.*, 2015).

These results highlight a clear resource role for quantum discord, specifically measured by operational quantifiers such as the IP and the DS, in black-box QEM settings, elucidating in particular how quantum correlations beyond entanglement manifest themselves as coherence in all local bases for the probing subsystem. Discordant states, i.e., all states but those of Eq. (14), are not only disturbed by all possible local measurements on A , but are also modified by — hence sensitive to — all nontrivial unitary evolutions on subsystem A . This is exactly the ingredient needed for the estimation and discrimination tasks described above.

In practice, however, one might want to assess the general purpose performance of probe states, rather than their worst case scenario only. One can then introduce alternative figures of merit quantifying how suitable a state is, on average, for estimation or discrimination of unitary

transformations, where the average is performed over all generators of a fixed spectral class. This can be done by replacing the minimum with an average according to the Haar measure, in Eqs. (17) and (25), respectively. Such a study has been carried out in (Farace *et al.*, 2016) by defining the local average Wigner-Yanase skew information, which corresponds to the average version of the DS in case the probing subsystem A is a qubit (Farace *et al.*, 2014).

Unlike the minimum, the average skew information is found not to be a measure of discord anymore. In particular, it vanishes only on states of the form $\frac{\mathbb{I}_A}{d_A} \otimes \tau^B$, that is, tensor product states between a maximally mixed state on A , and an arbitrary state on B (Farace *et al.*, 2016). This entails that, to ensure a reliable discrimination of local unitaries on average, the input states need to have either one of these two (typically competing) ingredients: nonzero local purity of the probing subsystem, or nonzero correlations (of any nature) with the ancilla. The interplay between the average performance and the minimum one, which instead relies on discord, as well as a study on the role of entanglement, are detailed in (Farace *et al.*, 2016).

Finally, we notice that the presence and use of quantum correlations beyond entanglement has also been investigated in other tasks related to metrology and illumination, such as ghost imaging with (unentangled) thermal source beams (Ragy and Adesso, 2012). Adopting a coarse-grained two-mode description of the beams, quantum discord was found to be relevant for the implementation of ghost imaging in the regime of low illumination, while more generally total correlations in the thermal source beams were shown to determine the quality of the imaging, as quantified by the signal-to-noise ratio.

III. IDENTICAL PARTICLES

Measuring devices and sensors operating with many-body systems are among the most promising instances where QEM can be actually experimented; indeed, their large numbers of elementary constituents play the role of resources according to which the accuracy of parameter estimation can be scaled. Typical instances in which the QEM paradigm has been studied are in fact interferometers based on ultracold atoms confined in optical lattices (Cronin *et al.*, 2009; Gerry and Knight, 2005; Giorgini *et al.*, 2008; Haroche and Raimond, 2006; Inguscio *et al.*, 2006; Köhl and Esslinger, 2006; Leggett, 2001, 2006; Pethick and Smith, 2004; Pitaevskii and Stringari, 2003; Yukalov, 2009) where a precise control on the state preparation and on the dynamics can nowadays be obtained. These systems are made of spatially confined bosons or fermions, *i.e.* of constituents behaving as identical particles, a fact that has not been properly taken into account in most of the

literature.

In systems of distinguishable particles, the notion of separability and entanglement is well-established (Horodecki *et al.*, 2009): it is strictly associated with the natural tensor product structure of the multi-particle Hilbert space and expresses the fact that one is able to identify each one of the constituent subsystems with their corresponding single-particle Hilbert spaces. On the contrary, in order to describe identical particles one must extract from the tensor product structure of the whole Hilbert space either the symmetric (bosonic) or the anti-symmetric (fermionic) sector (Feynman, 1994; Sakurai, 1994). This fact demands a more general approach to the notions of non-locality and entanglement based not on the particle aspect proper of first quantization, rather on the mode description typical of second quantization (Argentieri *et al.*, 2011; Barnum *et al.*, 2004, 2005; Benatti and Floreanini, 2014; Benatti *et al.*, 2010a, 2011, 2012a,b, 2014a,b; Marzolino, 2013; Narnhofer, 2004; Summers and Werner, 1985, 1987a,b; Zanardi *et al.*, 2004).

The notion of entanglement in many-body systems has already been addressed and discussed in the literature: for instance, see (Amico *et al.*, 2008; Balachandran *et al.*, 2013; Banuls *et al.*, 2006; Bloch *et al.*, 2008; Calabrese *et al.*, 2012; Dowling *et al.*, 2006; Eckert *et al.*, 2002; Ghirardi *et al.*, 2002; Grabowski *et al.*, 2011; Hines *et al.*, 2003; Kraus *et al.*, 2009; Lewenstein *et al.*, 2007; Li *et al.*, 2001; Micheli *et al.*, 2003; Modi *et al.*, 2012; Paskauskas and You, 2001; Schliemann *et al.*, 2001; Schuch *et al.*, 2004; Song *et al.*, 2012; Tichy *et al.*, 2013; Wiseman and Vaccaro, 2003; Y. Shi, 2004) and references therein. Nevertheless, only limited results actually apply to the case of identical particles and their applications to QEM. From the existing literature on possible metrological uses of identical particles, there emerges as a controversial issue the distinction between particle and mode entanglement. Before illustrating the general approach developed in (Benatti *et al.*, 2010a, 2012a, 2014a) within which this matter can be settled, we shortly overview the main aspects of the problem. Readers who feel that the discussion whether the states in question are to be considered as entangled or not is rather academic may be reassured by the very pragmatical result that independently of this discussion, systems of indistinguishable bosons offer a metrological advantage over distinguishable particles, in the sense that for certain measurements one would have to massively entangle the latter for obtaining the same sensitivity as one obtains “for free” from the symmetrized states of the former. This we show explicitly in Sec.III.B.

Entanglement based on the particle description proper of first quantization has been discussed for pure states in (Li *et al.*, 2001; Paskauskas and You, 2001;

Schliemann *et al.*, 2001). In the fermionic case, Slater determinants are identified as the only non-entangled fermionic pure states, for, due to the Pauli exclusion principle, they are the least correlated many-particle states. In the case of bosons, two inequivalent notions of particle entanglement have been put forward: in (Paskauskas and You, 2001), a pure bosonic state is declared non-entangled if and only if all particles are prepared in the same single particle state, thus leading to a bosonic product state. In (Li *et al.*, 2001), bosonic non-entangled pure states are identified as those corresponding to permanents (the bosonic analogue of Slater determinants); in other words, a pure bosonic state is non-entangled not only when, as in the first approach mentioned above, it is a product state, but also when all bosons are prepared in pairwise orthogonal single particle states. The particle-based entanglement was then studied in (Eckert *et al.*, 2002) both for bosons and fermions, and then generalized within a more mathematical and abstract setting in (Grabowski *et al.*, 2011, 2012).

From the point of view of particle description, a different perspective was provided in (Ghirardi *et al.*, 2002, 2004), based on the fact that non-entangled pure states should possess a complete set of local properties, identifiable by local measurements. For a more recent non standard approach, still based on the particle description, see (Lo Franco and Compagno, 2016).

In all these approaches, only bipartite systems consisting of two identical particles are discussed; however, in the fermionic case, simple necessary and sufficient criteria of many particle-entanglement, based on the single-particle reduced density matrix, have been elaborated in (Plastino *et al.*, 2009).

A change of perspective occurred in (Barnum *et al.*, 2004, 2005; Vedral, 2003; Zanardi *et al.*, 2004) where the focus moved from particles to orthonormal modes; within this approach, states that are not mode-entangled are characterized by correlations that can be explained in terms of joint classical occupation probabilities of the modes. It then follows that entanglement and non-locality depend on the mode description which has been chosen. In all cases, however, identical-particle states represented by fermionic Slater determinants or bosonic permanents can be neither particle nor mode-entangled, in direct conflict with claims that these states are particle-entangled and hence metrologically useful (Demkowicz-Dobrzański *et al.*, 2014). In an attempt to resolve the conflict, in (Killoran *et al.*, 2014) it is shown that such a pseudo, or “fluffy-bunny” entanglement as called in (Beenakker, 2006; Wiseman *et al.*, 2003), which is due to bosonic state symmetrization, can be turned into the entanglement of distinguishable, and thus metrologically accessible, modes; however, the operations needed for such a transformation are non-local in the mode picture and thus ultimately responsible for the achieved final non-local resources.

The variety of approaches regarding entanglement in identical particle systems and their use in metrological applications, in particular whether entanglement is necessary or not to achieve sub-shot-noise accuracies, can be looked at from the unifying point of view provided by the algebraic approach to quantum many-body systems proper of second quantization (Bratteli and Robinson, 1987; Strocchi, 1985). Within this scheme, we first address the relations between mode and particle entanglement. Then, we focus on specific quantum metrological issues and show that neither entangled states, nor preliminary state preparation as spin-squeezing are necessary in order to achieve sub shot-noise accuracies using systems of identical particles.

A. Particle and mode entanglement

In quantum mechanics, indistinguishable particles cannot be identified by specific labels, whence their states must be either completely symmetrized (bosons) or anti-symmetrized (fermions). This makes second quantization a most suited approach to deal with them, while the particle representation proper of first quantization results to be too restrictive for a consistent treatment of entangled identical particles.

As a simple example of the consequences of indistinguishability, consider two qubits, each of them described by the Hilbert space $\mathcal{H} = \mathbb{C}^2$, where we select the orthonormal basis $\{|0\rangle, |1\rangle\}$ of eigenvectors of σ_z : $\sigma_z|0\rangle = |0\rangle$, $\sigma_z|1\rangle = -|1\rangle$. If the two qubits describe distinguishable particles, the tensor product structure of the common Hilbert space $\mathbb{C}^4 = \mathbb{C}^2 \otimes \mathbb{C}^2$ exhibits the fact that one knows which is qubit 1 and which one qubit 2. Instead, if the two qubits are indistinguishable, the corresponding Hilbert space loses its tensor product structure and becomes, in the bosonic case, the three-dimensional subspace \mathbb{C}^3 spanned by the orthogonal symmetric vectors $|00\rangle$, $|11\rangle$ and $|+\rangle := (|01\rangle + |10\rangle)/\sqrt{2}$, whereas, in the Fermionic case, it reduces to the anti-symmetric vector $|-\rangle := (|01\rangle - |10\rangle)/\sqrt{2}$.

An important consequence of this fact is that generic mixed states of indistinguishable particles must be represented by density matrices that arise from convex linear combinations of projectors onto symmetrized or anti-symmetrized pure states. Indeed, while operators must be symmetrized in order to comply with particle indistinguishability, by symmetrizing density matrices of distinguishable particles, say sending $\rho = \rho_1 \otimes \rho_2$ into $(\rho_1 \otimes \rho_2 + \rho_2 \otimes \rho_1)/2$ one cannot in general obtain an appropriate bosonic or fermionic state, unlike sometimes stated in the literature (see e.g. (Demkowicz-Dobrzański *et al.*, 2014)). For instance, the two qubit symmetric mixed state $\rho \otimes \rho$ cannot be a fermionic density matrix since the only fermionic state is $|-\rangle\langle -|$. On the other hand, in the bosonic case, $\rho \otimes \rho$ can be a bona-fide bosonic state only if

ρ is pure. Indeed, as $\rho \otimes \rho|-\rangle = \det(\rho)|-\rangle$, if $\det(\rho) \neq 0$, then $\rho \otimes \rho$ has support also in the anti-symmetric component of the Hilbert space. A careful discussion of the relationship between permutationally invariant density matrices and symmetric/anti-symmetric states can be found in (Damanet *et al.*, 2016).

Instead of on the standard particle picture, second quantization is based on the so-called mode picture. In general, a mode is any of the normalized vectors $|\psi\rangle$ of the same single particle Hilbert space \mathcal{H} that is used to describe each one of a system of identical particles. In practice, one fixes an orthonormal basis $\{|\psi_i\rangle\}$ in \mathcal{H} and populates the i -th mode by acting on the so-called vacuum vector with (powers of) the creation operator a_i^\dagger , $a_i^\dagger|0\rangle = |\psi_i\rangle$, while the adjoint operators a_i annihilate the vacuum, $a_i|0\rangle = 0$. For bosons, $[a_i, a_j^\dagger] = a_i a_j^\dagger - a_j^\dagger a_i = \delta_{ij}$ and one can find arbitrarily many particles in a given mode, while for fermions $\{a_i, a_j^\dagger\} = a_i a_j^\dagger + a_j^\dagger a_i = \delta_{ij}$ and each mode can be occupied by one fermion, at most.

Typical modes are given by the eigenvectors of a given single particle Hamiltonian, orthogonal polarization directions, the left and right position in a double-well potential, or the atomic positions in an optical lattice. In the case of free photons, typical modes are plane waves labeled by wave vector and polarization, arising from the quantization of classical electrodynamics in terms of independent harmonic oscillators. Within this picture, saying that there are n photons in a certain energy mode can also be interpreted as a quantum oscillator being promoted to its n -th excited state.

In dealing with distinguishable particles, quantum entanglement (Horodecki *et al.*, 2009) is basically approached by referring to the tensor product structure of the total Hilbert space which embodies the fact that particles can be identified. From the previous discussion, it follows that, in the case of identical particles, one ought to consider the entanglement of modes rather than the entanglement of particles. The relations between the two approaches are studied in detail in (Benatti *et al.*, 2014b); here, we briefly compare them in the case of two two-mode indistinguishable bosons associated with orthogonal single-particle pure states $|\psi_{1,2}\rangle \in \mathcal{H}$, as, for instance, orthogonal polarization states, described by creation and annihilation operators a_i, a_i^\dagger , $i = 1, 2$. The pure state

$$|1, 1\rangle = a_1^\dagger a_2^\dagger |0\rangle \quad (26)$$

belongs to the two-particle sector of the symmetric Fock space and represents a pure state with one boson in each mode. In the particle picture of first quantization, the same state corresponds to the symmetrized vector

$$|\psi_+\rangle = \frac{1}{\sqrt{2}} \left(|\psi_1\rangle \otimes |\psi_2\rangle + |\psi_2\rangle \otimes |\psi_1\rangle \right), \quad (27)$$

which describes a balanced superposition of two states: one with the first particle in the state $|\psi_1\rangle$ and the second

particle in the state $|\psi_2\rangle$, the other with the two states exchanged. The net result is that, given the state (27), we can only say that one particle is surely in the state $|\psi_1\rangle$ and that the other one is surely in the state $|\psi_2\rangle$, but we cannot attribute a specific state to a specific particle.

As an example of the misunderstandings that can result from sticking to the particle picture when dealing with identical particles, state (27) is often referred to as entangled: actually, it is only formally so in the particle picture, while, as we shall show below, it is separable in the mode picture. Indeed, (27) would clearly display particle entanglement as the state cannot be written as a tensor product of single particle states. Instead, in (26), the same state is expressed as the action on the vacuum state of two independent creation operators and thus should correspond to a mode-separable state under whichever meaningful definition of mode-entanglement one should adopt. We shall later give solid ground to this latter statement, but firstly we show that, while for distinguishable particles the state (27) is the prototype of an entangled pure state, it is nevertheless separable for identical ones.

We follow the reformulation of separability of pure bipartite states as given in (Ghirardi *et al.*, 2002) where the starting point is that, whatever the mathematical characterization of entangled bipartite pure states, physically it must express the impossibility of attributing with certainty specific properties to any of the parties.

Given a one-particle state $|\varphi\rangle \in \mathcal{H}$, the projection $P_\varphi = |\varphi\rangle\langle\varphi|$ corresponds to one particle being in the state $|\varphi\rangle$. This property is possessed by the system if its state ρ is such that $\text{Tr}(\rho P_\varphi) = 1$, that is if and only if $\rho = P_\varphi$. For distinguishable particles, separable pure states are tensor products of single particle pure states, while local operators address each particle independently. Such are, for instance, the projectors $P_1 = |\psi_1\rangle\langle\psi_1| \otimes \mathbb{I}$, $P_2 = \mathbb{I} \otimes |\psi_2\rangle\langle\psi_2|$ and $P_3 = P_1 P_2$, where \mathbb{I} denotes the identity and $|\psi_{1,2}\rangle \in \mathcal{H}$. If the two-particle system is in the state $|\psi\rangle \equiv |\psi_1\rangle \otimes |\psi_2\rangle$, the projectors P_1 , P_2 and P_3 correspond to the possessed properties “first particle in the state $|\psi_1\rangle$ independently of the second particle”, “second particle in the state $|\psi_2\rangle$ independently of the first one” and “first particle in the state $|\psi_1\rangle$ and second particle in the state $|\psi_2\rangle$ ”. Indeed, $|\psi\rangle$ is eigenstate of $P_{1,2}$ and of P_3 so that $\langle\psi|P_1|\psi\rangle = \langle\psi|P_2|\psi\rangle = \langle\psi|P_3|\psi\rangle = 1$.

For distinguishable particles, the state $|\psi_+\rangle$ in (27) is such that $P_{1,2,3}$ do not correspond to possessed properties; indeed, $\langle\psi_+|P_i|\psi_+\rangle = \frac{1}{2}$, $i = 1, 2, 3$. Actually, $P_{1,2,3}$ describe possessed properties only for separable states $|\Psi\rangle = |\psi_1\rangle \otimes |\psi_2\rangle$. In other terms, the entangled pure states of two distinguishable particles are such that no individual properties can be attributed to each one or both of the constituent parties.

Suppose now the two particles to be identical. As already observed, this implies that their pure states must be symmetric for bosons and antisymmetric for fermions

under particle exchange. Note that particle identity does not exclude that individual properties might be attributed to the two parties; the only constraint is that no property can be attributed to a definite party otherwise it could be used to distinguish it from the other one. Consider the single-particle property described by the one-dimensional projector $|\psi\rangle\langle\psi|$; for two indistinguishable particles the property “one particle is in the state $|\psi\rangle$ ”, must then be represented by the symmetric projection

$$\mathcal{E}_\psi = |\psi\rangle\langle\psi| \left(\mathbb{I} - |\psi\rangle\langle\psi| \right) + \left(\mathbb{I} - |\psi\rangle\langle\psi| \right) \otimes |\psi\rangle\langle\psi| + |\psi\rangle\langle\psi| \otimes |\psi\rangle\langle\psi|.$$

Similarly, the projector corresponding to the two qubits possessing two different properties $|\psi_{1,2}\rangle\langle\psi_{1,2}|$ must be $P_{12}^{symm} = |\psi_1\rangle\langle\psi_1| \otimes |\psi_2\rangle\langle\psi_2| + |\psi_2\rangle\langle\psi_2| \otimes |\psi_1\rangle\langle\psi_1|$. This is a projection if and only if $\langle\psi_1|\psi_2\rangle = 0$, in which case it assumes the form $P_{12}^{symm} = \mathcal{E}_{\psi_1} \mathcal{E}_{\psi_2}$. Therefore, despite the formal entanglement of $|\psi_+\rangle$, properties can be attributed to both its parties. Indeed,

$$\langle\psi_+|\mathcal{E}_{\psi_1}|\psi_+\rangle = \langle\psi_+|\mathcal{E}_{\psi_2}|\psi_+\rangle = \langle\psi_+|P_{12}^{symm}|\psi_+\rangle = 1.$$

Then, as entanglement is associated with the impossibility of attributing individual properties, despite its form, the state in (27) is still separable¹.

The above approach to identical particle entanglement has been developed within a first quantization context, focusing on single particle states. From the point of view of second quantization, separability and entanglement are instead to be related to the algebraic structure of Bose and Fermi systems rather than to the possibility of attributing individual properties: this is the point of view usually adopted in the analysis of many-body systems, where the primary object of investigation are the algebras of operators rather than their representations on particular Hilbert spaces (Bratteli and Robinson, 1987; Emch, 1972; Haag, 1992; Strocchi, 1985, 2008a,b, 2012; Thirring, 2002).

In order to appropriately formulate the notion of entanglement in systems made of identical particles, the leading intuition is that there is no a priori given tensor product structure, reminiscent of particle identification, either in the Hilbert space or in the algebra of observables. Therefore, questions about entanglement and separability are meaningful only with reference to specific classes of observables. Within this broader context, entanglement becomes a caption for the non-local

¹ If $|\psi_+\rangle$ is the state of two distinguishable particles, it remains an eigenstate of $\mathcal{E}_{P_{1,2}}$; however, these observables do not represent complete sets of properties for a single particle. Indeed, these complete sets of properties are associated to the pure state projectors $|\psi_{1,2}\rangle\langle\psi_{1,2}|$.

quantum correlations between observables exhibited by certain quantum states (Clifton and Halvorson, 2001; Halvorson and Clifton, 2000; Keyl *et al.*, 2006, 2003; Moriya, 2006; Summers and Werner, 1985, 1987a,b; Verch and Werner, 2005).

Polynomials in creation and annihilation operators can be used to generate bosonic and fermionic algebras \mathcal{A} containing all physically relevant many-body observables. Quite in general, physical states are given by positive and normalized linear functionals $\omega : \mathcal{A} \rightarrow \mathbb{C}$ associating to any operator $A \in \mathcal{A}$ its mean value $\omega(A)$, such that positive observables $A \geq 0$ are mapped to positive numbers $\omega(A) \geq 0$ and $\omega(\mathbb{I}) = 1$. Typical instances of states are the standard expectations obtained by tracing with respect to a given density matrix ρ , $\omega(A) = \text{Tr}(\rho A)$; notice, however, that in presence of infinitely many degrees of freedom, not all physically meaningful states, like, for instance, the thermal ones, can be represented by density matrices (Narnhofer, 2004).

Within such an algebraic approach, one can study the entanglement between observables of \mathcal{A} with respect to a state ω by considering *algebraic bipartitions* (Benatti *et al.*, 2010a, 2014a), namely any couple of subalgebras $(\mathcal{A}_1, \mathcal{A}_2)$ of \mathcal{A} , having only the identity in common, $\mathcal{A}_1 \cap \mathcal{A}_2 = \mathbb{I}$, such that the linear span of products of their operators generate the entire algebra, $\mathcal{A}_1 \cup \mathcal{A}_2 = \mathcal{A}$ and

1. $[A_1, A_2] = 0$ for all $A_1 \in \mathcal{A}_1$ and $A_2 \in \mathcal{A}_2$ in the bosonic case;
2. $[A_1^e, A_2] = 0$ for all even elements $A_1^e \in \mathcal{A}_1^e$ and for all $A_2 \in \mathcal{A}_2$, and, similarly, $[A_1, A_2^e] = 0$ for all even elements $A_2^e \in \mathcal{A}_2^e$ and for all $A_1 \in \mathcal{A}_1$, in the fermionic case,

where the even elements $A_{1,2}^e \in \mathcal{A}_{1,2}$, are those remaining invariant under the transformation of creation and annihilation operators $\vartheta(a_i^\dagger) = -a_i^\dagger$ and similarly for a_i . They generate the so-called even subalgebras $\mathcal{A}_{1,2}^e \subset \mathcal{A}$ which contain the only fermionic operators accessible to experiments.

These requests embody the notion of algebraic independence and generalize the tensor product structure typical of the particle picture. Within this more general context, an operator \mathcal{O} is called $(\mathcal{A}_1, \mathcal{A}_2)$ -local if of the form $\mathcal{O} = A_1 A_2$, $A_1 \in \mathcal{A}_1$ and $A_2 \in \mathcal{A}_2$. Furthermore, a state ω is called $(\mathcal{A}_1, \mathcal{A}_2)$ -separable if the expectation values of all $(\mathcal{A}_1, \mathcal{A}_2)$ -local operators can be decomposed into convex combinations of expectations:

$$\omega(A_1 A_2) = \sum_i \lambda_k \omega_k^{(1)}(A_1) \omega_k^{(2)}(A_2), \quad (28)$$

in terms of other states $\omega_i^{(1)}, \omega_i^{(2)}$ with $\lambda_k > 0$, $\sum_i \lambda_k = 1$; otherwise, ω is called $(\mathcal{A}_1, \mathcal{A}_2)$ -entangled².

² This generalized definition of separability can be easily extended

In the standard case of bipartite entanglement for pairs of distinguishable particles (Horodecki *et al.*, 2009), entangled states are all density matrices ρ which cannot be written in the form

$$\rho = \sum_k \lambda_k \rho_k^{(1)} \otimes \rho_k^{(2)}, \quad \lambda_k \geq 0, \quad \sum_k \lambda_k = 1,$$

with $\rho_k^{(1)}$ and $\rho_k^{(2)}$ density matrices of the two parties. That the algebraic definition reduces to the standard one becomes apparent in the case of two qubits by choosing the algebraic bipartition $\mathcal{A}_1 = \mathcal{A}_2 = \mathcal{M}_2$, where \mathcal{M}_2 is the algebra of 2×2 complex matrices and the expectation value $\omega(A_1 A_2) = \text{Tr}(\rho A_1 \otimes A_2)$.

Henceforth, we shall focus on many-body systems, whose elementary constituents can be found in M different states or modes described by a discrete set of annihilation and creation operators $\{a_i, a_i^\dagger\}_{i \in I}$; this is a very general framework, useful for the description of physical systems in quantum optics, in atomic and condensed matter physics. A bipartition of the M -mode algebra \mathcal{A} associated with the system can then be easily obtained by considering two disjoint sets $\{a_i, a_i^\dagger | i = 1, 2, \dots, m\}$ and $\{a_j, a_j^\dagger | j = m+1, m+2, \dots, M\}$, M being possibly infinite. The two sets form subalgebras $\mathcal{A}_{1,2}$ that indeed constitute an algebraic bipartition of \mathcal{A} ; in practice, it is determined by the integer $0 < m < M$.

As for distinguishable particles, the case of pure states is easier. Pure states in the algebraic context³ correspond to those expectations on \mathcal{A} that cannot be written as convex combinations of other expectations. For them the separability condition (28) simplifies: one can indeed prove (Benatti and Floreanini, 2016) that pure states ω on \mathcal{A} are separable with respect to a given bipartition $(\mathcal{A}_1, \mathcal{A}_2)$ if and only if

$$\omega(A_1 A_2) = \omega(A_1) \omega(A_2), \quad (29)$$

for all local operators $A_1 A_2$. In other terms, separable pure states are just product states.

B. Mode entanglement and metrology: bosons

The differences between mode-entanglement and standard entanglement are best appreciated in the case of N bosons that can occupy M different modes; this is a very general situation encountered, for instance,

to the case of more than two partitions; for instance, in the case of an n -partition, Eq.(28) would extend to $\omega(A_1 A_2 \dots A_n) = \sum_k \lambda_k \omega_k^{(1)}(A_1) \omega_k^{(2)}(A_2) \dots \omega_k^{(n)}(A_n)$.

³ In the algebraic description, Hilbert spaces are a byproduct of the algebraic structure and of the expectation functional (state) defined on it (Bratteli and Robinson, 1987).

in ultracold gases consisting of bosonic atoms confined in a multiple site optical lattices. These systems turn out to be a unique laboratory for studying quantum effects in many-body physics, *e.g.* in quantum phase transition and matter interference phenomena, and also for applications in quantum information (*e.g.* see (Cronin *et al.*, 2009; Gerry and Knight, 2005; Giorgini *et al.*, 2008; Haroche and Raimond, 2006; Inguscio *et al.*, 2006; Köhl and Esslinger, 2006; Leggett, 2006; Pethick and Smith, 2004; Pitaevskii and Stringari, 2003; Yukalov, 2009), and references therein).

The algebra \mathcal{A} is in this case generated by creation and annihilation operators $a_i^\dagger, a_i, i = 1, 2, \dots, M$, obeying the commutation relations, $[a_i, a_j^\dagger] = \delta_{ij}$. The reference state ω is given by the expectations with respect to the vacuum state $|0\rangle$, $a_i|0\rangle = 0$ for all $1 \leq i \leq M$, so that the natural states whose entanglement properties need be studied are vectors in the Fock Hilbert space \mathcal{H} spanned by the many-body Fock states

$$|n_1, n_2, \dots, n_M\rangle = \frac{1}{\sqrt{n_1! n_2! \dots n_M!}} (a_1^\dagger)^{n_1} (a_2^\dagger)^{n_2} \dots (a_M^\dagger)^{n_M} |0\rangle, \quad (30)$$

or density matrices acting on it.

Given a bipartition $(\mathcal{A}_1, \mathcal{A}_2)$ defined by two disjoint groups of creation and annihilation operators, any element $A_1 \in \mathcal{A}_1$, commutes with any element $A_2 \in \mathcal{A}_2$, *i.e.* $[A_1, A_2] = 0$.

In this framework, a necessary and sufficient condition for pure states $|\psi\rangle$ to be separable with respect to a given bipartition $(\mathcal{A}_1, \mathcal{A}_2)$ is that they be generated by acting on the vacuum state with $(\mathcal{A}_1, \mathcal{A}_2)$ -local operators (Benatti *et al.*, 2012a),

$$|\psi\rangle = \mathcal{P}(a_1^\dagger, \dots, a_m^\dagger) \cdot \mathcal{Q}(a_{m+1}^\dagger, \dots, a_M^\dagger) |0\rangle, \quad (31)$$

where \mathcal{P}, \mathcal{Q} are polynomials in the creation operators relative to the first m modes, the last $M - m$ modes, respectively. Pure states that can not be cast in the above form are thus $(\mathcal{A}_1, \mathcal{A}_2)$ -entangled.

When the state of the bosonic many-body system is not pure, it can be described by a density matrix ρ , that, using the Fock basis (30), can be written as:

$$\rho = \sum_{\{n\}, \{n'\}} \rho_{\{n\}, \{n'\}} |n_1, n_2, \dots, n_M\rangle \langle n'_1, n'_2, \dots, n'_M|, \quad (32)$$

$$\sum_{\{n\}} \rho_{\{n\}, \{n\}} = 1.$$

Since density matrices form a convex set whose extremals are given by the projectors onto pure states, one deduces that mixed states ρ as in (32) are $(\mathcal{A}_1, \mathcal{A}_2)$ -separable if and only if they are convex combinations of $(\mathcal{A}_1, \mathcal{A}_2)$ -separable one-dimensional projectors states.

An interesting application of these general considerations is given by a system of N bosons that can be found

in just two modes, $M = 2$. In the Bose-Hubbard approximation, N ultracold bosonic atoms confined in a double-well potential can be effectively described in this way. The two creation operators a_1^\dagger and a_2^\dagger generate out of the vacuum bosons in the two wells, so that the Fock basis (30) can be conveniently relabeled in terms of the integer k counting the number of bosons in the first well:

$$|k, N - k\rangle = \frac{(a_1^\dagger)^k (a_2^\dagger)^{N-k}}{\sqrt{k!(N-k)!}} |0\rangle, \quad 0 \leq k \leq N. \quad (33)$$

Furthermore, a_1, a_1^\dagger , respectively a_2, a_2^\dagger , generate two commuting subalgebras \mathcal{A}_1 and \mathcal{A}_2 that, together, in turn generate the whole algebra \mathcal{A} ; it is the only nontrivial bipartition of the system.

Then, the states $|k, N - k\rangle$ are separable: this agrees with the fact that they are created by the local operators $(a_1^\dagger)^k (a_2^\dagger)^{N-k}$. Indeed, for any polynomial operator $\mathcal{P}_1 \in \mathcal{A}_1$ and $\mathcal{P}_2 \in \mathcal{A}_2$, the expectation value of the product $\mathcal{P}_1 \mathcal{P}_2$ is such that (compare with (29))

$$\begin{aligned} \langle k, N - k | \mathcal{P}_1 \mathcal{P}_2 | k, N - k \rangle \\ = \frac{1}{k!(N-k)!} \langle 0 | a_1^k \mathcal{P}_1 (a_1^\dagger)^k | 0 \rangle \langle 0 | a_2^{N-k} \mathcal{P}_2 (a_2^\dagger)^{N-k} | 0 \rangle \\ = \langle k | \mathcal{P}_1 | k \rangle \langle N - k | \mathcal{P}_2 | N - k \rangle, \end{aligned}$$

where $|k\rangle$ and $|N - k\rangle$ are single-mode Fock states. Consequently, mixed separable states must be diagonal with respect to the Fock basis (33), *i.e.* density matrices of the form (Benatti *et al.*, 2010a):

$$\rho = \sum_{k=0}^N p_k |k, N - k\rangle \langle k, N - k|, \quad p_k \geq 0, \quad \sum_{k=0}^N p_k = 1. \quad (34)$$

Unlike $\mathcal{P}_1 \mathcal{P}_2$, most observables of physical interest are non-local with respect to the bipartition $(\mathcal{A}_1, \mathcal{A}_2)$. Prominent among them are those used in phase estimation protocols based on ultra-cold atoms trapped in double-well interferometers. These are the collective bilinear operators:

$$\begin{aligned} J_x &= \frac{1}{2} (a_1^\dagger a_2 + a_1 a_2^\dagger), \\ J_y &= \frac{1}{2i} (a_1^\dagger a_2 - a_1 a_2^\dagger), \\ J_z &= \frac{1}{2} (a_1^\dagger a_1 - a_2^\dagger a_2), \end{aligned} \quad (35)$$

satisfying the $su(2)$ algebraic relations $[J_i, J_j] = i\epsilon_{ijk} J_k$, $i, j, k = x, y, z$. Notice that, although the operators in (35), as well as the exponentials $e^{i\theta J_x}$ and $e^{i\theta J_y}$, are non-local with respect to the bipartition $(\mathcal{A}_1, \mathcal{A}_2)$, $\theta \in [0, 2\pi]$, the exponential of J_z turns out to be local:

$$e^{i\theta J_z} = e^{i\theta a_1^\dagger a_1 / 2} e^{-i\theta a_2^\dagger a_2 / 2}, \quad (36)$$

with $e^{i\theta a_1^\dagger a_1 / 2} \in \mathcal{A}_1$ and $e^{-i\theta a_2^\dagger a_2 / 2} \in \mathcal{A}_2$. By a linear transformation, one can always pass to new annihilation

operators

$$b_1 = \frac{a_1 + a_2}{\sqrt{2}}, \quad b_2 = \frac{a_1 - a_2}{\sqrt{2}}, \quad (37)$$

and corresponding creation operators $b_{1,2}^\dagger$, and rewrite the three operators in (35) as

$$\begin{aligned} J_x &= \frac{1}{2}(b_1^\dagger b_1 - b_2^\dagger b_2), \\ J_y &= \frac{1}{2i}(b_1 b_2^\dagger - b_1^\dagger b_2), \\ J_z &= \frac{1}{2}(b_1 b_2^\dagger + b_1^\dagger b_2). \end{aligned} \quad (38)$$

To b_i, b_i^\dagger , $i = 1, 2$ one associates the bipartition $(\mathcal{B}_1, \mathcal{B}_2)$ of \mathcal{A} consisting of the subalgebras generated by b_1, b_1^\dagger , respectively b_2, b_2^\dagger : with respect to it, the exponential of J_x becomes:

$$e^{i\theta J_x} = e^{i\theta b_1^\dagger b_1/2} e^{-i\theta b_2^\dagger b_2/2}, \quad (39)$$

with $e^{i\theta b_1^\dagger b_1/2} \in \mathcal{B}_1$ and $e^{-i\theta b_2^\dagger b_2/2} \in \mathcal{B}_2$. This explicitly shows that an operator, local with respect to a given bipartition, can result non-local if a different algebraic bipartition is chosen.

These considerations are relevant for metrological applications (Boixo *et al.*, 2009; Bollinger *et al.*, 1996; Bouyer and Kasevich, 1997; Caves, 1981; Dorner *et al.*, 2009; Dowling, 1998, 2008; Dunningham *et al.*, 2002; Giovannetti *et al.*, 2004, 2011; Higgins *et al.*, 2007; Holland *et al.*, 2002; Holland and Burnett, 1993; Kacprowicz *et al.*, 2010; Kitagawa and Ueda, 1993; Korbicz *et al.*, 2005; Pezzè and Smerzi, 2009; Sanders and Milburn, 1995; Sørensen *et al.*, 2001; Tóth *et al.*, 2009; Uys and Meystre, 2007; Wang and Sanders, 2003; Wineland *et al.*, 1994; Yurke, 1986; Yurke *et al.*, 1986). In fact, ultracold atoms trapped in a double-well optical potential realize a very accurate interferometric device: state preparation and beam splitting can be precisely achieved by tuning the interatomic interaction and by acting on the height of the potential barrier. The combination of standard Mach-Zehnder (MZ) type interferometric operations, *i.e.* state preparation, beam splitting, phase shift and subsequent beam recombination, can be effectively described as a suitable rotation of the initial state ρ_{in} by a unitary transformation (Sanders and Milburn, 1995; Yurke *et al.*, 1986):

$$\rho_{\text{in}} \mapsto \rho_\theta = U_\theta \rho_{\text{in}} U_\theta^\dagger, \quad U_\theta = e^{i\theta J_n}. \quad (40)$$

The phase change is induced precisely by the operators in (35) through the combination:

$$J_n \equiv n_x J_x + n_y J_y + n_z J_z, \quad n_x^2 + n_y^2 + n_z^2 = 1. \quad (41)$$

In practice, the state transformation $\rho_{\text{in}} \mapsto \rho_\theta$ inside the interferometer can be effectively modeled as a pseudo-spin rotation along the unit vector $\mathbf{n} = (n_x, n_y, n_z)$,

whose choice depends on the specific realization of the interferometric apparatus and of the adopted measurement procedure.

As discussed in Section I.A, in the case of distinguishable particles, for any separable state ρ_{sep} the quantum Fisher information is bounded by $I_\theta[\rho_{\text{in}}, J_n] \leq cN$ where c is a constant independent of N that can be taken as the maximum quantum Fisher information of a single system in all components of the mixed separable state (see eq.(9)) (Fujiwara and Hashizumé, 2002; Giovannetti *et al.*, 2006; Pezzè and Smerzi, 2009). Then, entangled initial states of distinguishable particles are needed in order to obtain sub-shot-noise accuracies. A corresponding statement was proven in (Benatti and Braun, 2013) for a system of N indistinguishable bosons: The quantum Fisher information with respect to a parameter θ imprinted onto a $(\mathcal{A}, \mathcal{B})$ -separable state via a $(\mathcal{A}, \mathcal{B})$ -local unitary operator of the form $U(\theta) = \exp(i(A(a, a^\dagger) + B(b, b^\dagger)))$, where A (respectively B) are hermitian functions of a, a^\dagger (respectively b, b^\dagger), strictly vanishes. Hence, in order to be able to estimate θ at all (and even more so to beat the SQL in terms of the total number of bosons), the separability of the input state or the locality of the unitary operator that imprints θ on the state need to be broken. Fortunately, mode non-locality is easily achieved *e.g.* in quantum optics with a simple beam-splitter, without any particle interactions (see below).

In the case of the double-well system introduced above, the operator algebras defined by a_1^\dagger, a_1 and a_2^\dagger, a_2 , are the natural ones. With respect to them, a balanced Fock number state as in (33) of the form $\rho_{N/2} = |\frac{N}{2}, \frac{N}{2}\rangle\langle\frac{N}{2}, \frac{N}{2}|$ is separable. By choosing the vector \mathbf{n} in the plane orthogonal to the z direction one computes $I_\theta[\rho_{N/2}, J_n] = \frac{N^2}{2} + N$, thus approaching the Heisenberg limit.

As a concrete example (Benatti and Braun, 2013), consider the action of a beam-splitter described by the unitary $U_{BS}(\alpha) = \exp(\alpha a_1^\dagger a_2 - \alpha^* a_1 a_2^\dagger)$ involving two modes. If the complex transparency parameter $\alpha = i\theta/2$, with θ real, then $U_{BS}(\alpha) = \exp(i\theta J_x)$. In the case of N distinguishable qubits, a state as the balanced Fock number state $|\frac{N}{2}, \frac{N}{2}\rangle$ has half qubits in the state $|0\rangle$ and half in the state $|1\rangle$ such that $\sigma_z|i\rangle = (-)^i|i\rangle$, $i = 0, 1$. Then, $J_i = \sum_{j=1}^N \frac{\sigma_i^{(j)}}{2}$ with $i \in \{x, y, z\}$ have zero mean values, while the purity of $\rho_{N/2}$ yields $I_\theta[\rho_N, J_x] = 4\langle J_x^2 \rangle = N$, since

$$J_x^2 = \frac{N}{4} + \frac{1}{4} \sum_{j \neq k=1}^N \sigma_x^{(j)} \sigma_x^{(k)}.$$

Instead, in the case of indistinguishable bosonic qubits, the mean values of J_i for all $i \in \{x, y, z\}$ in (35) vanish,

while using (35), $I_\theta[\rho_N, J_x] = 4 \langle J_x^2 \rangle = \frac{N^2}{2} + N$.

Unlike in the case of distinguishable particles, the quantum Fisher information can thus attain a value greater than N even with initial states like $\rho_{N/2}$ that are separable with respect to a given bipartition. As mentioned before, the rotation operated by the beam-splitter is not around the z axis and is thus non-local with respect to the chosen bipartition. It is thus not the entanglement of the states fed into the beam-splitter that helps overcoming the shot-noise-limit in the transparency parameter θ estimation accuracy; rather, the non-local character of the rotations operated by the apparatus on initially separable states allows $\sigma(\theta_{\text{est}})$ to be smaller than $1/\sqrt{N}$, with the possibility of eventually reaching the Heisenberg $1/N$ limit (Benatti and Braun, 2013; Benatti *et al.*, 2011).

Of course, by changing bipartition $(\mathcal{B}_1, \mathcal{B}_2)$ via equations (37), the action of the beam-splitter, as outlined in (39), is local, exactly as in the pseudo-spin formalism. In this bipartition, the non-locality necessary for enhancing the sensitivity completely resides in the state. Therefore, the mode-description, leaves the freedom to locate the resources necessary to accuracy-enhancing either in the entanglement of the state or in the non-locality of the operations performed on it.

In (Benatti and Braun, 2013), also the paradigmatic case of phase estimation in a Mach-Zehnder interferometer was considered with similar results: in the mode-bipartition corresponding to the two modes after the first beam-splitter, the phase shift operation in one arm is a local operator. Hence, at that stage the state must be mode-entangled to allow estimating the phase shift with an accuracy better than the shot-noise limit. However, under certain conditions the first beam-splitter can generate enough mode-entanglement from a mode-separable state fed into the two input ports of the interferometer to beat the SQL. For more general settings, the question of what scaling of the quantum Fisher information can be achieved with the number of indistinguishable bosons is still open. Notice that non-locality is partially attributed to operations like beam-splitting, rather than entirely to states also in the different context of transforming “fluffy-bunny” entanglement into useful entanglement developed in (Killoran *et al.*, 2014).

For massive bosons one might think that the tensor product of Fock states is a very natural state: As the boson-number is conserved at the energies considered, one cannot make coherent superpositions of different numbers of atoms. However, in typical experiments, one needs to average over many runs, and the real difficulty consists in controlling the atom number with single-atom precision from run to run. One has

therefore effectively a mixed state with a distribution of different atom numbers. So far the best experimentally demonstrated approximations to Fock states with massive bosons are number-squeezed states with a few dB of squeezing (Esteve *et al.*, 2008; Gross *et al.*, 2010b; Riedel *et al.*, 2010b). One may hope that novel measurement techniques such as the quantum gas microscope (Bakr *et al.*, 2009) may enable precise knowledge of boson numbers in the future and thus the preparation of Fock states with a large number of atoms at least in a post-selected fashion.

Recently, the authors of (Oszmaniec *et al.*, 2016) used concentration of measure techniques to investigate the usefulness of randomly sampled probe states for unitary quantum metrology. They show that random pure states drawn from the Hilbert space of distinguishable particles typically do not lead to super-classical scaling of precision. However, random states from the symmetric subspace, i.e. bosonic states, typically achieve the HL, even for very mixed isospectral states. Moreover, the quantum enhancement is typically robust against the loss of particles, in contrast to e.g. GHZ-states. It remains to be seen how entangled these states typically are in the sense of mode-entanglement, but independently of the outcome of such an assessment, these results are in line with the finding that the naturally symmetrized pure states of bosons are a useful resource for quantum metrology.

C. Mode entanglement and metrology: fermions

As in the case of bosonic systems, we shall consider generic fermionic many-body systems made of N elementary constituents that can occupy M different states or modes, $N < M$. The creation a_i^\dagger and annihilation a_i operators for mode i obey now the anticommutation relations $\{a_i, a_j^\dagger\} = \delta_{ij}$ and generate the fermion algebra \mathcal{A} , i.e. the norm closure of all polynomials in these operators. As already specified before, a bipartition $(\mathcal{A}_1, \mathcal{A}_2)$ of this algebra is the splitting of the collection of creation and annihilation operators into two disjoint sets. The Hilbert space \mathcal{H} of the system is again generated out of the vacuum state $|0\rangle$ by the action of the creation operators; it is spanned by the many-body Fock states,

$$|n_1, n_2, \dots, n_M\rangle = (a_1^\dagger)^{n_1} (a_2^\dagger)^{n_2} \dots (a_M^\dagger)^{n_M} |0\rangle, \quad (42)$$

where the integers n_1, n_2, \dots, n_M are the occupation numbers of the different modes, with $\sum_i n_i = N$; they can now take only the two values 0 or 1.

As already clear from the definition of fermionic algebraic bipartitions, because of the anti-commutation relations, in dealing with fermions, one must distinguish between even and odd operators. While the even component \mathcal{A}^e of \mathcal{A} consists of elements $A^e \in \mathcal{A}$ such that $\vartheta(A^e) = A^e$, the odd component \mathcal{A}^o of \mathcal{A} consists of

those elements $A^o \in \mathcal{A}$ such that $\vartheta(A^o) = -A^o$. Even elements of \mathcal{A} commute with all other elements, while odd elements commute only with even ones.

The anticommuting character of the fermion algebra \mathcal{A} puts stringent constraints on the form of the fermion states that can be represented as product of other states, like the ones appearing in the decomposition (28) that defines separable states. Specifically, as a consequence of the result in (Araki and Moriya, 2003), any product $\omega_k^{(1)}(A_1) \omega_k^{(2)}(A_2)$ vanishes whenever A_1 and A_2 both belong to the odd components of their respective subalgebras. Then, given a mode bipartition $(\mathcal{A}_1, \mathcal{A}_2)$ of the fermionic algebra \mathcal{A} , *i.e.* a decomposition of \mathcal{A} in the subalgebra \mathcal{A}_1 generated by the first m modes and the subalgebra \mathcal{A}_2 , generated by the remaining $M - m$ ones, it follows that the decomposition (28) is meaningful only for local operators $A_1 A_2$ for which $[A_1, A_2] = 0$, so that the definition of separability it encodes is completely equivalent to the one adopted for bosonic systems.

As a further consequence of the result in (Araki and Moriya, 2003), one derives that if a state ω is non vanishing on a local operator $A_1^o A_2^o$, with the two components $A_1^o \in \mathcal{A}_1^o$, $A_2^o \in \mathcal{A}_2^o$ both belonging to the odd part of the two subalgebras, then ω is entangled with respect to the bipartition $(\mathcal{A}_1, \mathcal{A}_2)$. Indeed, if $\omega(A_1^o A_2^o) \neq 0$, then ω cannot be written as in (28), and therefore cannot be separable.

Using these results, as for the bosonic case, one shows that (Benatti *et al.*, 2014a), given a bipartition of the fermionic algebra \mathcal{A} determined by the integer m , a pure state $|\psi\rangle$ results separable if and only if it can be written in the form (31). Examples of pure separable states of N fermions are the Fock states in (42); indeed, they can be recast in the form

$$|k_1, \dots, k_m; p_{m+1}, \dots, p_M\rangle = \left[(\hat{a}_1^\dagger)^{k_1} \dots (\hat{a}_m^\dagger)^{k_m} \right] \times \left[(\hat{a}_{m+1}^\dagger)^{p_{m+1}} \dots (\hat{a}_M^\dagger)^{p_M} \right] |0\rangle,$$

where the \mathcal{P} and \mathcal{Q} appearing in (31) are now monomials in the creation operators of the two partitions.

Concerning the metrological use of fermionic systems, the situation may appear more problematic than with bosons, for each mode can accommodate at most one fermion; therefore, the scaling with N of the sensitivity in the estimation of physical parameters may worsen. Indeed, while a two-mode bosonic apparatus, as a double-well interferometer, filled with N particles is sufficient to reach sub shot-noise sensitivities, with fermions, a multimode interferometer is needed in order to reach comparable sensitivities (for the use of multi-mode interferometers see, for instance (Cooper *et al.*, 2009; D'Ariano *et al.*, 1998; D'Ariano and Paris, 1997; Söderholm *et al.*, 2003; Vourdas and Dunningham, 2005), and (Cooper *et al.*, 2012) in the fermionic case). As an example, consider a system of N fermions in M modes, with M even, and

let us fix the balanced bipartition $(M/2, M/2)$, in which each of the two parts contain $m = M/2$ modes, taking for simplicity $N \leq m$. As generator of the unitary transformation $\rho \rightarrow \rho_\theta$ inside the measuring apparatus let us take the following operator:

$$\mathcal{J} = \frac{1}{2} \sum_{k=1}^m \omega_k \left(a_k^\dagger a_{m+k} + a_{m+k}^\dagger a_k \right), \quad (43)$$

where ω_k is a given spectral function, *e.g.* $\omega_k \simeq k^p$, with p positive. The apparatus implementing the above state transformation is clearly non-local with respect to the chosen bipartition: $e^{i\theta\mathcal{J}}$ can not be written as the product $A_1 A_2$ of two components made of operators referring only to the first, second partition, respectively. It represents a generalized, multimode beam splitter, and the whole measuring device behaves as a multimode interferometer.

Let us feed the interferometer with a pure initial state, $\rho = |\psi\rangle\langle\psi|$,

$$|\psi\rangle = |\underbrace{1, \dots, 1}_N, \underbrace{0, \dots, 0}_{m-N}; \underbrace{0, \dots, 0}_m\rangle = a_1^\dagger a_2^\dagger \dots a_N^\dagger |0\rangle, \quad (44)$$

where the fermions occupy the first N modes of the first partition; $|\psi\rangle$ is a Fock state and therefore it is separable, as already discussed. The quantum Fisher information can be easily computed (Benatti *et al.*, 2014a)

$$I_\theta[\rho, \mathcal{J}] = \sum_{k=1}^N \omega_k^2. \quad (45)$$

Unless ω_k is k -independent, $I_\theta[\rho, \mathcal{J}]$ is larger than N and therefore the interferometric apparatus can beat the shot-noise limit in θ -estimation, even starting with a separable state. Actually, for $\omega_k \simeq k^p$, one gets: $I_\theta[\rho, \mathcal{J}] \simeq N^{2p+1}$, reaching sub-Heisenberg sensitivities with a linear device. Note that this result and the ability to go beyond the Heisenberg limit is not a “geometrical” phenomenon attributable to a phase accumulation even on empty modes (D'Ariano and Paris, 1997); rather, it is a genuine quantum effect, that scales as a function of the number of fermions, the resource available in the measure.

Again, as in the bosonic case, it is not the entanglement of the initial state that helps overcoming the shot-noise-limit in the phase estimation; rather, it is the non-local character of the rotations operated by the apparatus on an initially separable state that allows one beating the shot-noise limit.

IV. MORE GENERAL HAMILTONIANS

The assumption of a hamiltonian $H = \sum_{k=1}^N h_k$ considered for the derivation of eq.(2), where Λ and λ are the largest and smallest eigenvalues of h_k , respectively, not

only implies distinguishable subsystems, it is also restrictive in two other important regards: a.) The existence of such bounds on the spectrum of h_k may not be warranted, and b.) interactions between the subsystems are excluded. In this section we explore the consequences of lifting these restrictions.

A. Lifting spectral limitations

A large portion of the work on QEM stems from quantum optics, where the basic dynamical objects are modes of the e.m. field, corresponding to simple harmonic oscillators, $h_k = \hbar\omega_k(n_k + 1/2)$. A phase shift in mode k can be implemented by $U = \exp(in_k\theta)$. Clearly, for the relevant hamiltonian $h_k = n_k$ acting as generator of the phase shift, $\Lambda = \infty$ and $\lambda = 0$. Hence, eq.(2) would imply a minimal uncertainty $\text{Var}(\theta_{\text{est}}) = 0$. Of course, one may argue that in reality one can never use states of infinite energy, such that there is effectively a maximum energy. However, it need not be that the maximum energy sustainable by the system must be distributed over N modes. Indeed, what is typically counted in quantum optics in terms of resources is not the number of modes N , but the total number of photons n , directly linked to the total energy. It turns out, that the total number of modes (or subsystems, in general) is completely irrelevant for achieving optimal sensitivity, even if the parameter is coded in several modes or subsystems, e.g. with a general unitary transformation of the form $U = \exp(i\theta \sum_{k=1}^N h_k)$, if one can stock the same amount of energy in a single system as in the total system. Note that this is often the case in quantum optics, where different modes can be spatially confined or parametrically influenced by the same optical elements such as mirrors, beam-splitters, and phase shifters whose material properties ultimately determine the maximum amount of energy that can be used.

To see the liberating effect of unbound spectra, recall that for any initial pure state $|\psi\rangle$ propagated by a hamiltonian of the form $H = \theta G$ with a hermitian generator G for a time T the QFI is given by (Braunstein and Caves, 1994; Braunstein *et al.*, 1996)

$$I_\theta = 4\text{Var}(G)T^2 \equiv 4(\langle G^2 \rangle - \langle G \rangle^2)T^2. \quad (46)$$

Let $G = \sum_i e_i |i\rangle\langle i|$ be the spectral decomposition of G , and $|\psi\rangle = \sum_{i=1}^L c_i |i\rangle$, where we assume that $|1\rangle$ ($|L\rangle$) are the states of lowest (largest) energy available. Then $\text{Var}(G) = \sum_{i=1}^L p_i e_i^2 - (\sum_{i=1}^L p_i e_i)^2$ with $p_i = |c_i|^2$ and $\sum_{i=1}^L p_i = 1$. The Popoviciu inequality (Popoviciu, 1935) states $\text{Var}(G) \leq (e_L - e_1)^2/4$. It is saturated for $p_1 = p_L = 1/2$, $p_i = 0$ else. The state $|\psi\rangle = (|1\rangle + e^{i\varphi}|L\rangle)/\sqrt{2}$ with an arbitrary phase φ saturates the inequality and thus maximizes I_θ . If e_L

or e_1 is degenerate, only the total probability for the degenerate energy levels is fixed to $1/2$, and arbitrary linear combinations in the degenerate subspace are allowed. But the value of $\text{Var}(G)$ remains unchanged under such redistributions, and we may still choose just two non-vanishing probabilities $p_1 = p_L = 1/2$. The derivation did not make use of a multi-mode structure of the energy eigenstates. Hence, exactly the same minimal uncertainty of θ_{est} can be obtained by superposing the ground state of a single mode with a Fock state of given maximum allowed energy as with an arbitrarily entangled multi-mode state containing components of up to the same maximum energy.

For a specific example, consider phase estimation in a Mach-Zehnder interferometer. It has $N = 2$ modes, and a phase shift just in one of them, i.e. the relevant hamiltonian is $H = \theta n_1$. Adding energy conservation of the two modes at the beam-splitters (i.e. the fact that the accessible states are two-mode Fock states of the form $|n - n_2, n_2\rangle$, where n_2 with $0 \leq n_2 \leq n$ is the number of photons in the second mode), one immediately finds that the optimal two mode state is $(|n, 0\rangle + |0, n\rangle)/\sqrt{2}$, i.e. the highly-entangled N00N state (Boto *et al.*, 2000). However, we can achieve exactly the same variance of G and hence sensitivity with the single-mode state $(|n\rangle + |0\rangle)/\sqrt{2} \otimes \rho_2$, i.e. a product state where we keep the second mode in *any* state ρ_2 . In both cases the maximum energy of the first mode is $n\hbar\omega$ (assuming $\omega_1 = \omega_2 = \omega$), and the average energy in the interferometer $n\hbar\omega/2$ (neglecting the vacuum energy $\hbar\omega/2$). Hence, also from the perspective of maximum energy deposit in the optical components, there is no advantage in using two entangled modes. If the Mach-Zehnder interferometer is realized abstractly via Ramsey-pulses on N two-level systems (states $|0\rangle, |1\rangle$) for the beam-splitters, and a phase shift $\exp(i\theta J_z)$, $J_z = \sum_{i=1}^N \sigma_z^{(i)}/2$, the state that maximizes $\text{var}(J_z)$ is the (maximally entangled) GHZ state $(|0\dots 0\rangle + |1\dots 1\rangle)/\sqrt{2}$. But exactly the same uncertainty can be obtained with a single spin- j ($j = N/2$) in the state $(|j, j\rangle + |j, -j\rangle)/\sqrt{2}$ (in the usual $|j, m\rangle$ notation, where j is the total angular momentum and m its z -component). Clearly, allowing as large a spectrum for a single system as for the combined systems makes entanglement entirely unnecessary here.

These considerations teach us that the relevant quantity to be maximized is the quantum uncertainty of the generator G . This can be understood in terms of generalized Heisenberg uncertainty relations, where the generator G plays the quantity complementary to θ , as was found early on (Braunstein *et al.*, 1996). In a multi-component system maximizing $\text{Var}(G)$ may be achieved with highly entangled states, but if the spectral range of a single system admits the same $\text{Var}(G)$, there is no need for entanglement.

In (Braun, 2011, 2012) a state of the form $(|0\rangle + |2n\rangle)/\sqrt{2}$ was found to be optimal for mass measurements with a nano-mechanical oscillator given a maximum allowed number of excitation quanta $2n$ and times much larger than the oscillation period (for shorter times there are contributions also from the dependence on frequency of the energy eigenfunctions). The same state of a single mode of the e.m. field is optimal for measuring the speed of light (Braun *et al.*, 2015). In both cases the quantum uncertainties scale as $1/n$ (QFI proportional to n^2), and obviously no entanglement is needed. Of course, a state of the form $(|0\rangle + |2n\rangle)/\sqrt{2}$ (called “half a N00N” state in (Braun, 2011)) is still highly non-classical. In fact, already a single Fock state $|n\rangle$ is highly non-classical, as is witnessed by its highly oscillatory Wigner function (Schleich, 2001) with substantial negative parts. The superposition $(|0\rangle + |2n\rangle)/\sqrt{2}$ leads in addition to $2n$ lobes in azimuthal direction that explain the sensitivity of phase measurements $\propto n^{-1}$. Alternatively one can use coherent state superpositions (“CSS states”) (Bimbarb *et al.*, 2010; Lund *et al.*, 2004; Neergaard-Nielsen *et al.*, 2006; Suzuki *et al.*, 2006; Wakui *et al.*, 2007; Yukawa *et al.*, 2013), i.e. “Schrödinger-cat” type states of the form $(|\alpha\rangle + |-\alpha\rangle)/\sqrt{2}$. They have been created in quantum optics with values of $\alpha = 0.79$ in (Ourjoumtsev *et al.*, 2006). In (Lund *et al.*, 2004) a “breeding method” based on weak squeezing, beam mixing with an auxiliary coherent field, and photon detecting with threshold detectors was proposed to achieve values up to $\alpha \leq 2.5$, but the success probability has been found to be too low for a realistic iterated protocol. An alternative based on homodyning was proposed in (Etesse *et al.*, 2014; Laghaout *et al.*, 2013) and implemented in (Etesse *et al.*, 2015), leading to $\alpha \simeq 1.63$. The current record in the optical domain for “large” α appears to be $\alpha \simeq \sqrt{3}$, achieved from two-mode squeezed vacuum and n -photon detection on one of the modes (Huang *et al.*, 2015). In the microwave regime, CSS states with $\alpha = \sqrt{7}$ have been generated, as well as superposition of coherent states with smaller phase differences with up to 111 photons (Vlastakis *et al.*, 2013). Almost arbitrary superpositions with a small number of photons can be generated by using couplings of a mode with two-level systems that can be tuned in and out of resonance, but reviewing the entire literature of non-classical states in general is beyond the scope of the present review (see e.g. (Deléglise *et al.*, 2008; Gottesman *et al.*, 2001; Hofheinz *et al.*, 2009) and the Nobel lectures of Serge Haroche and David J. Wineland (Haroche, 2013; Wineland, 2013) for historical accounts of the development of these fields and many more references).

B. Non-linear hamiltonians

Most discussions of QEM consider, implicitly or explicitly, evolution under a hamiltonian that is linear in a collective variable of the system. For illustration, consider Ramsey spectroscopy on a collection of N atoms with ground and excited states $|g\rangle$, $|e\rangle$, respectively. If these have energies $\pm\hbar\omega/2$, then the hamiltonian giving rise to the Ramsey oscillation is

$$H = \hbar\omega S_z \equiv \hbar\omega \sum_{i=1}^N s_z^{(i)} \quad (47)$$

where $s_z \equiv \frac{1}{2}(|e\rangle\langle e| - |g\rangle\langle g|)$ is a pseudo spin-1/2 operator describing the transition and the superscript (i) indicates the i -th atom. This hamiltonian is manifestly linear in S_z , and can be trivially decomposed into microhamiltonians $H = \sum_i h^{(i)} \equiv \sum_i \hbar\omega s_z^{(i)}$ that describe the uncoupled precession of each atom in the ensemble. For a single Ramsey sequence of duration T , so that the unknown parameter is $\theta = \omega T$, the SQL and HL sensitivities are as described in Section I.A, $\text{Var}(\theta_{\text{est}})_{\text{SQL}} = N^{-1}$ and $\text{Var}(\theta_{\text{est}})_{\text{HL}} = N^{-2}$. The assumption of uncoupled particles is often physically reasonable, for example when describing photons in a linear interferometer or low-density atomic gases for which collisional interactions can be neglected. In other systems, including Bose-Einstein condensates (Gross *et al.*, 2010a; Riedel *et al.*, 2010a), laser interferometers at high power (Aasi *et al.*, 2013; The LIGO Scientific Collaboration, 2011), high-density atomic magnetometers (Dang *et al.*, 2010; Kominis *et al.*, 2003; Shah *et al.*, 2010; Vasilakis *et al.*, 2011) and high-density atomic clocks (Deutsch *et al.*, 2010), the assumption of uncoupled particles is unwarranted. This motivates the study of nonlinear hamiltonians.

The unusual features of nonlinear hamiltonians are well illustrated in the following example (Boixo *et al.*, 2008b). First, define the collective operator $S_0 \equiv \sum_i s_0^{(i)} \equiv \sum_i \mathbb{1}^{(i)}$, where $\mathbb{1}$ indicates the identity operator. S_0 is clearly the total number of particles. Now consider the nonlinear hamiltonian

$$H = \hbar\Omega S_0 S_z = \hbar\Omega \sum_{i=1}^N \sum_{j=1}^N s_0^{(i)} s_z^{(j)}. \quad (48)$$

This is linear in the unknown Ω , but of second order in the collective variables S . At the microscopic level, the hamiltonian describes a pair-wise interaction, with energy $\hbar\Omega s_0^{(i)} s_z^{(j)}$, between each pair of particles (i, j) .

For a system with a fixed number N of particles, the consequence for the dynamics of the system is very simple: the operator S_0 can be replaced by its eigenvalue N , leading to an effective hamiltonian

$$H_N = \hbar N \Omega S_z. \quad (49)$$

Estimation of the product $N\Omega T = N\Theta$ now gives the same uncertainties that we saw earlier in the estimation of θ . That is, $\text{Var}(N\Theta_{\text{est}})_{\text{SQL}} = N^{-1}$ so that $\text{Var}(\Theta_{\text{est}})_{\text{SQL}} = N^{-3}$. Similarly, $\text{Var}(\Theta_{\text{est}})_{\text{HL}} = N^{-4}$. More generally, a nonlinear hamiltonian containing k -order products of collective variables will contain N^k microscopic interaction terms that contribute to the hamiltonian and thus to the rate of change of an observable such as S_z under time evolution. In contrast, the variance of such a macroscopic observable, e.g. $\text{Var}(S_z)$, scales as N^1 (SQL) or N^0 (HL). This allows signal-to-noise ratios scaling as N^{2k-1} (SQL) and as N^{2k} (HL) (Boixo *et al.*, 2007). These noise terms depend only on the size of the system and nature of the initial state but not on the hamiltonian (Boixo *et al.*, 2007).

That nonlinearities lead to improved scaling of the sensitivity appears to have been independently discovered by A. Luis and J. Beltrán (Beltrán and Luis, 2005; Luis, 2004, 2007) and by S. Boixo and co-workers (Boixo *et al.*, 2008a,b, 2007; Datta and Shaji, 2012). A related proposal using interactions to give a scaling $\sigma(\theta_{\text{est}}) \propto 2^{-N}$ is described in (Roy and Braunstein, 2008).

Prior to the appearance of these results, the term “Heisenberg limit,” which was introduced into the literature by (Holland and Burnett, 1993) in the context of interferometric phase estimation with the definition $\sigma(\phi_{\text{est}}) = 1/N$, had been used, often indiscriminately, to describe 1) the sensitivity $\sigma(\phi_{\text{est}}) = 1/N$, 2) the scaling $\sigma(\phi_{\text{est}}) \propto 1/N$, 3) the best possible sensitivity with N particles, and 4) the best possible scaling with N particles (Giovannetti *et al.*, 2004, 2006, 2011). Clearly these multiple definitions are not all compatible in a scenario with a nonlinear hamiltonian. Taking as a definition “error... bounded by the inverse of the physical resources,” and implicitly considering scaling, (Zwierz *et al.*, 2010) (see also (Zwierz *et al.*, 2011)) showed that an appropriate definition for “physical resource” is the query complexity of the system viewed as a quantum network.

For the simplest optical nonlinearity, $\theta \propto N^2$ and quadrature detection, it has been shown that quadrature squeezed states are near-optimal (Maldonado-Mundo and Luis, 2009). Considering the same nonlinearity and limiting to classical inputs, i.e. coherent states and mixtures thereof, it is argued in (Rivas and Luis, 2010) that non-linear strategies can out-perform linear ones by concentrating the available particles in a small number of high-intensity probes. (Tilma *et al.*, 2010) analyzed a variety of entangled coherent states for nonlinear interferometry of varying orders, and found that in most cases entanglement degraded the sensitivity for high-order nonlinearities. (Berrada, 2013) considered the use of two-mode squeezed states as inputs to a non-linear interferometer, including the effects of loss, and showed a robust advantage for such states.

As already mentioned in Section I.B, the above re-

sults concern local measures implying some prior knowledge of the parameter being estimated. The situation for global measures without prior information is considered in (Hall and Wiseman, 2012).

C. Proposed experimental realizations

A number of physical systems have been proposed for nonlinear QEM: Propagation through nonlinear optical materials (Beltrán and Luis, 2005; Luis, 2004, 2007), collisional interactions in Bose-Einstein condensates (Boixo *et al.*, 2009), Duffing nonlinearity in nanomechanical resonators (Woolley *et al.*, 2008), and nonlinear Faraday rotation probing of an atomic ensemble (Napolitano and Mitchell, 2010).

1. Nonlinear optics

The first proposals concerned nonlinear optics (Beltrán and Luis, 2005; Luis, 2004, 2007), in which a nonlinear optical susceptibility is directly responsible for a phase-shift $\theta \propto N^k$, where k is the order of the nonlinear contribution to the refractive index. In the simplest example, (Beltrán and Luis, 2005) showed that an input coherent state $|\alpha\rangle$, experiencing a Kerr-type nonlinearity described by the unitary $\exp[i\Theta(a^\dagger a)^2]$, and detected in quadrature $X = \frac{1}{\sqrt{2}}(a + a^\dagger)$, gives an outcome distribution

$$P(X = x|\Theta) = |\langle x|e^{i\Theta(a^\dagger a)^2}|\alpha\rangle|^2. \quad (50)$$

If we consider the case of small Θ , imaginary $\alpha = i\sqrt{\langle N \rangle}$, and the estimator $\Theta_{\text{est}} = \bar{X}/|\partial\bar{X}/\partial\theta|$, where $\bar{X} \equiv \sum_{i=1}^M X_i$ is the mean of the observed quadratures, we find the standard deviation

$$\sigma(\Theta_{\text{est}}) = \frac{\sigma(X)}{\sqrt{M}|\partial\bar{X}/\partial\theta|} = \frac{\sigma(X)}{\sqrt{M}|\langle [X, (a^\dagger a)^2] \rangle|} \quad (51)$$

$$= \frac{1}{4M^{1/2}N^{-3/2}}. \quad (52)$$

Here we have used $\sigma(X) = 1/\sqrt{2}$ for the quantum mechanical uncertainty of X in the initial state, and which up to corrections of order Θ^2 holds also for the evolved state when Θ is small. For large N , this strategy saturates the QCRB; the QFI is straightforwardly calculated to be $I_\Theta = 4N + 24N^2 + 16N^3$.

2. Ultra-cold atoms

Coherent interaction-based processes are well developed in Bose-Einstein condensates and have been used extensively for squeezing generation. For example, a confined two-species Bose condensate experiences collisional

energy shifts described by an effective hamiltonian

$$H_{\text{eff}} \propto a_{11}n_1(n_1 - 1) + 2a_{12}n_1n_2 + a_{22}n_2(n_2 - 1) \quad (53)$$

$$= (a_{11} - 2a_{12} + a_{22})S_z^2 + 2(a_{11} - a_{22})S_zS_0$$

$$+ \text{terms in } S_z, S_0 \quad (54)$$

where $S_z \equiv \frac{1}{2}(n_1 - n_2)$ and $S_0 \equiv \frac{1}{2}(n_1 + n_2)$ are pseudo-spin operators, n_1 and n_2 are the number of atoms of species 1 and 2, respectively, and a_{ij} are the collisional scattering lengths. In ^{87}Rb , and with $|1\rangle \equiv |F=1, m=-1\rangle$, $|2\rangle \equiv |F=2, m=1\rangle$, the scattering lengths (near zero magnetic field) have the ratio $a_{11} : a_{12} : a_{22} = 1.03 : 1 : 0.97$. A proven method to generate spin squeezing in this system is to increase a_{12} using a Feshbach resonance to give the single-axis twisting hamiltonian $H_{\text{eff}} \propto S_z^2$ (Muessel *et al.*, 2014), plus terms proportional to S_z and S_0 , which induce a global rotation and a global phase shift, respectively.

It was observed in (Boixo *et al.*, 2008a) that the zero-field scattering lengths naturally give $a_{11} - 2a_{12} + a_{22} \approx 0$, making small the coefficient of S_z^2 and leaving the S_0S_z term as the dominant nonlinear contribution. Detailed analyses of the BEC physics beyond the simplified single-mode treatment here are given in (Boixo *et al.*, 2009, 2008b; Tacla and Caves, 2013). The strategy gives $N^{-3/2}$ scaling for estimation of the relative scattering length $a_{11} - a_{22}$. (Mahmud *et al.*, 2014) describe a strategy of dynamical decoupling to suppress the second-order terms in the hamiltonian and thus make dominant three-body interactions, giving a sensitivity scaling of $N^{-5/2}$ for measurements of three-body collision strengths.

3. Nano-mechanical oscillators

(Woolley *et al.*, 2008) propose a nonlinear interferometer using two modes of a nano-mechanical oscillator, with amplitudes x_a and x_b , experiencing the nonlinear Hamiltonian

$$H_{\text{eff}} = H_{\text{SHO}}^{(a)} + H_{\text{SHO}}^{(b)} + \frac{1}{4}\chi_a m\omega^2 x_a^4 + \frac{1}{4}\chi_b m\omega^2 x_b^4 + C(t) \quad (55)$$

where H_{SHO} indicates the simple harmonic oscillator hamiltonian, χ is the Duffing nonlinearity coefficient, ω is the low-amplitude resonance frequency, and C is an externally-controlled coupling between modes a and b that produces a beam-splitter interaction. With an interferometric sequence resembling a Mach-Zehnder interferometer, the Duffing nonlinearity can be estimated with uncertainty scaling as $N^{-3/2}$, where N is the number of excitations.

4. Nonlinear Faraday rotation

Whereas Luis and co-workers considered phenomenological models of optical nonlinearities,

(Napolitano and Mitchell, 2010) describes an *ab-initio* calculation of the optical nonlinearity produced on a particular atomic transition, using degenerate perturbation theory and a collective quantum variable description. This gives an effective Hamiltonian for the interaction of polarized light, described by the Stokes operators \mathbf{S} , with the collective orientation and alignment spin variables \mathbf{J} of an atomic ensemble:

$$H_{\text{eff}} = H_{\text{eff}}^{(2)} + H_{\text{eff}}^{(4)} + O(\mathbf{S}^3) \quad (56)$$

$$H_{\text{eff}}^{(2)} = \alpha^{(1)} S_z J_z + \alpha^{(2)} (S_x J_x + S_y J_y) \quad (57)$$

$$H_{\text{eff}}^{(4)} = \beta_J^{(0)} S_z^2 J_0 + \beta_N^{(0)} S_z^2 N_A + \beta^{(1)} S_0 S_z J_z$$

$$+ \beta^{(2)} S_0 (S_x J_x + S_y J_y). \quad (58)$$

where the α and β coefficients are linear and non-linear polarizabilities that depend on the detuning of the probe photons from the atomic resonance. By proper choice of detuning and initial atomic polarization \mathbf{J} , the term $\beta^{(1)} S_0 S_z J_z$ can be made dominant, making the hamiltonian formally equivalent to that of Eq. (48). Note that $\beta^{(1)} J_z$, proportional to the atomic polarization J_z , plays the role of the unknown interaction energy $\hbar\Omega$. The photons are thus made to interact, mediated by and proportional to the atomic polarization J_z . For a different detuning, the term $\alpha^{(1)} S_z J_z$ becomes dominant, allowing a linear measurement of the same quantity J_z with the same atomic system.

The experimental realization using a cold, optically-trapped ^{87}Rb atomic ensemble is described in (Napolitano *et al.*, 2011). The experiment observed the predicted scaling of $\text{Var}(J_z) \propto N^{-3}$ over a range of photon numbers from $N = 5 \times 10^5$ to 5×10^7 . For larger photon numbers the scaling worsened, i.e. $\text{Var}(J_z(N))$ had a logarithmic derivative > -3 . Due to this limited range of the N^{-3} scaling, and the difference in pre-factors $\beta^{(1)}$ versus $\alpha^{(1)}$, the nonlinear estimation never surpassed the sensitivity $\text{Var}(J_z)$ of the linear measurement for the same number of photons. Nonetheless, due to a shorter measurement time τ , the nonlinear measurement did surpass the linear measurement in spectral noise density $\text{Var}(J_z)\tau$, a common figure of merit for time- or frequency-resolved measurements.

D. Observations and commentary

Several differences between linear and nonlinear strategies, perhaps surprising, deserve comment.

First, it should be obvious that there is no conflict with the Heisenberg uncertainty principle (HUP). θ and Θ are parameters, not observables, and as such are not subject to operator-based uncertainty relations, neither the HUP nor generalizations such as the Robertson uncertainty relation (Robertson, 1929). Moreover, the advantageous scaling in $\delta\Theta$ is the result of a rapidly-growing *signal*,

rather than a rapidly decreasing statistical noise. A nonlinear hamiltonian immediately leads to a strong change in the scaling of the signal: even the simplest $k = 2$ non-linearity gives signal growing as $\omega \propto N^2$ and thus SQL uncertainty $(\delta\Theta)_{\text{SQL}} \propto N^{-3/2}$, which scales faster with N than does $(\delta\theta)_{\text{HL}} \propto N^{-1}$.

Second, the estimated phases θ and Θ necessarily reflect different physical quantities. $\hbar\omega$ describes a single-particle energy such as that due to an external field, whereas $\hbar\Omega$ describes a pair-wise interaction energy. As such, the uncertainties $\delta\theta$ and $\delta\Theta$ are not *directly* comparable. Any comparison of the efficacy of the measurements must introduce another element, a connection between a third physical quantity, θ , and Θ . This we have seen in Section IV.C.4, where the unknown J_z appears in both the linear and nonlinear hamiltonians. In what follows, we describe an optimized linear/nonlinear comparison.

E. Nonlinear measurement under number-optimized conditions

A more extensive study of nonlinear spin measurements, using the same system as (Napolitano *et al.*, 2011) is reported in (Sewell *et al.*, 2014). This work compared two estimation strategies, one linear and one non-linear, for measuring the collective variable J_y , which describes a component of the spin alignment tensor used in a style of optical magnetometry known as alignment-to-orientation conversion (Budker *et al.*, 2000; Pustelny *et al.*, 2008; Sewell *et al.*, 2012). The linear estimation used the term $\alpha^{(2)}S_yJ_y$, which appears in Eq. (57) and produces a rotation from linearly polarized light toward elliptically-polarized light. The nonlinear estimation in contrast used Eq. (57) in second order: in the first step, due to the $\alpha^{(2)}S_xJ_x$ term and the input S_x optical polarization, an initial J_y atomic polarization is rotated toward J_z by an angle $\phi \propto \langle S_x \rangle$, and thus $\propto N$, where N is the number of photons. In the second step, the term $\alpha^{(1)}S_zJ_z$ produces a Faraday rotation, i.e. from S_x toward S_y , by an angle proportional to the J_z polarization produced in the first step. The resulting S_y polarization is $S_y \propto J_yN^2$, while the statistical noise is $\sigma(S_y) \propto N^{1/2}$, giving sensitivity scaling $\sigma(J_y) \propto N^{-3/2}$. Importantly, the two estimation strategies use the same S_x -polarized input, and thus have identical statistical noise and cause identical damage in the form of spontaneous scattering, which adds noise to the atomic polarization.

The experimentally-observed nonlinear sensitivity was compared against the calculated ideal sensitivity of the linear measurement. Owing to its faster scaling, and more favorable pre-factors than in (Napolitano *et al.*, 2011), the nonlinear measurement's sensitivity surpassed that of the ideal linear measurement at about 2×10^7 photons. A comparison was also made when each mea-

surement was independently optimized for number N and detuning, which affects both the pre-factors α and the scattering. The ability to produce measurement-induced spin-squeezing was taken as the figure of merit, and the fully-optimized nonlinear measurement gave more squeezing than the fully-optimized linear measurement. This shows that for some quantities of practical interest, a nonlinear measurement can out-perform the best possible linear measurement.

F. Signal amplification with nonlinear hamiltonians

The single-axis twisting hamiltonian $H_{\text{twist}} = \chi S_z^2$, in addition to producing spin squeezed states, has been proposed as a nonlinear amplifier to facilitate state readout in atom interferometry (Davis *et al.*, 2016). Starting from an x -polarized coherent spin state $|\mathbf{x}\rangle$, and defining the unitary $U \equiv \exp[-iH_{\text{twist}}\tau/\hbar]$, the Wigner distribution of the squeezed state $U|\mathbf{x}\rangle$ is thin in the z direction, and is thus sensitive to rotations $\mathcal{R}_y(\phi)$ about the y -axis, so that states of the form $\mathcal{R}_y(\phi)U|\mathbf{x}\rangle$ have large quantum Fisher information with respect to ϕ . Exploitation of this in-principle sensitivity is challenging, however, because it requires low-noise readout, detecting S_z at the single-atom level if the Heisenberg limit is to be approached. In contrast, a sequence that applies H_{twist} , waits for rotation about the y -axis and then applies $-H_{\text{twist}}$ for an equal time generates the state $U^\dagger \mathcal{R}_y(\phi)U|\mathbf{x}\rangle$. Because U^\dagger is unitary, the quantum Fisher information is unchanged, but the perturbation implied by $\mathcal{R}_y(\phi)$ now manifests itself at the scale of the original coherent spin state, which is to say it is amplified from the HL scale up to the SQL scale, greatly facilitating detection. While this strategy clearly uses entanglement, it is nonetheless striking that un-doing the entanglement-generation step provides an important benefit.

G. Decoherence-enhanced measurements

Decoherence is arguably the most fundamental issue that plagues quantum-enhancements of all kinds, and quantum-enhanced measurements are no exception. However, decoherence has interesting physical properties which imply that it can also be useful for precision measurements. This goes beyond the benefits of decoherence and open system dynamics found as early as the late 1990s and the early 2000s, when it was realized that entanglement can be created through decay processes or more generally through coupling to common environments (Benatti *et al.*, 2009, 2010b, 2003, 2008; Braun, 2002, 2005; Plenio *et al.*, 1999), and, paradoxically, that decoherence of quantum computations can be reduced by rather strong dissipation that confines the computation to a decoherence-free subspace (DFS) through a

Zeno-type effect due to the rapid decay of states outside the DFS (Beige *et al.*, 2000a,b). Recently such ideas have found renewed interest, and meanwhile techniques have been proposed to create steady state entanglement in driven open quantum systems, such as cold Rydberg gases in the Rydberg-blockade regime (Lee *et al.*, 2015). It remains to be seen whether such stabilized entangled states are useful for precision measurements.

Here, on the contrary, we focus on the dynamics created by decoherence processes themselves. Decoherence arises from an interaction with an environment described by a non-trivial interaction hamiltonian $H_{\text{int}} = \sum_i S_i B_i$ with a similar structure as the non-linear hamiltonians considered above, where, however, one distinguishes operators S_i pertaining to the system and others (B_i) pertaining to the environment. The environment is typically considered as heat bath with a large number of degrees of freedom that may not be entirely accessible (Benatti and Floreanini, 2005; Breuer and Petruccione, 2002; Weiss, 1999). In addition, system and heat bath have their own hamiltonian, such that the total hamiltonian reads $H = H_S + H_B + H_{\text{int}}$. From simple model systems it is known that decoherence tends to become extremely fast for quantum superpositions of states that differ macroscopically in terms of the eigenvalues of one of the S_i . For example it was shown that superposing two Gaussian wave-packages of a free particle in one dimension that is coupled through its position to a heat bath of harmonic oscillators leads to decoherence times that scale as powers of \hbar that depend on how the wave packages are localized in phase space: The shortest decoherence time, scaling as $\hbar/|q_1 - q_2|$ results from wave packages distinguished only by positions q_1, q_2 , the longest one $\sim \hbar^{1/2}/\sqrt{|p_1 - p_2|}$ from wave packages that differ only in their mean momenta p_1, p_2 , and an intermediate one scaling as $\hbar^{2/3}/|(q_1 - q_2)(p_1 - p_2)|^{1/3}$. For systems of finite Hilbert-space dimensions such as angular momenta, one can often identify an effective \hbar that scales like the inverse Hilbert-space dimension which suggests that monitoring the decoherence process can lead to highly sensitive measurements, possibly surpassing the $1/\sqrt{N}$ scaling of the SQL.

That this intuition is correct was shown in (Braun and Martin, 2011), where a method was proposed for measuring the length of a cavity by monitoring the decoherence process of N atoms inside the cavity. The atoms are initially prepared in a highly excited dark state, where destructive interference prevents them from transferring their energy to a mode of the cavity with which they are resonant. For example, if one has two atoms coupling via an interaction $(g_1 \sigma_-^{(1)} + g_2 \sigma_-^{(2)})^\dagger + h.c.$ to a mode of the cavity with annihilation operator a , a state $\propto (1/g_1)|10\rangle - (1/g_2)|01\rangle$ of the atoms (where $|0\rangle$ and $|1\rangle$ are the ground and excited states of the atoms) is a dark state, also known as decoherence-free state: the amplitudes of photon transfer from the two atoms

to the cavity cancel. However, the couplings g_i depend on the position of the atoms relative to the cavity due to the envelope of the e.m. field. If the cavity changes its length L with the atoms at fixed positions, the g_i change, such that the original state becomes slightly bright. This manifests itself through the transfer of atomic excitations to the mode of the cavity, from where photons can escape and be detected outside. In (Braun and Martin, 2011) it was shown that through this procedure the minimal uncertainty with which L can be estimated (quantum Cramér-Rao bound) scales as $1/N$ even when using an initial product state of $N/2$ pairs of atoms. This scaling applies both for a perfect cavity and in the bad cavity limit, where superradiance arises. Thus, at least in principle, Heisenberg-limited scaling can arise here without the need for an entangled state, and in spite of the inherently decoherent nature of superradiance. However, the prefactor matters also here: Superradiance leads to a rapid decay of all states that are not dark, such that the available signal and with it the prefactor of the $1/N$ scaling law deteriorate rapidly with time.

Given the delocalized nature of the cavity mode in this example, it is possible in principle to make the number of atoms that interact with the mode arbitrarily large, in contrast to the non-linear schemes above, where the interactions have to decrease if the total energy is to remain an extensive quantity. But if the volume is kept fixed the atoms will start to interact so that the simple independent atom model of superradiance breaks down. For even larger N one has to increase the size of the cavity in order to accommodate all atoms. When the largest possible density is reached, the volume will have to grow $\propto N$, which implies coupling constants of the atoms to the cavity that decay as $1/\sqrt{N}$ and leads back to SQL scaling. In addition, the number of atoms has to be macroscopic in order to compete with the best classical sensitivities reached with interferometers such as LIGO: Assuming that the prefactor in the $1/N$ -scaling is of order one, one needs $\sim 10^{21}$ atoms for a minimal uncertainty of 10^{-21} . A cubic optical lattice with one atom every μm , the lattice would have to span 10m in x, y, z -direction in order to accommodate that many atoms. When using a diamond with regularly arranged NV-centers every 10 nm (such dense samples have been fabricated (Acosta *et al.*, 2009)), one would still require a cubic diamond of edge length 10cm. These examples show that competing with the best classical techniques is very challenging even if one can achieve HL scaling, as in classical protocols it is relatively easy to scale up the number of photons, compensating thus for a less favorable scaling with N .

H. Coherent averaging

There is nothing inherently quantum about the $1/\sqrt{N}$ scaling of the SQL. Rather, this behavior is a simple consequence of the central limit theorem applied to N independently acquired measurement results that are averaged as part of a classical noise-reduction procedure. The idea of coherent averaging is to replace the averaging by a coherent procedure, where the N probes all interact with a central quantum system (a “quantum bus”). In the end one measures the quantum bus or the entire system.

A simple example shows how this can lead to HL sensitivity without needing any entanglement: Consider N spins-1/2 interacting with a central spin-1/2 with the Ising-interaction $H_{\text{int}} = \sum_{i=1}^N g_i \sigma_z^{(0)} \otimes \sigma_z^{(i)}$, where the index zero indicates the central spin. The interaction commutes with the free hamiltonian of all spins $H_s = \hbar \sum_{i=0}^N \omega_i \sigma_z^{(i)}$, and we can solve the time-evolution exactly, starting from the initial product state $|\psi(0)\rangle = \frac{1}{\sqrt{2}}(|0\rangle_0 + |1\rangle_0) \otimes_{i=1}^N |0\rangle_i$. At time t , the state has evolved to $|\psi(t)\rangle = \frac{1}{\sqrt{2}}(e^{i(\omega_0/2 - N\bar{g})t}|0\rangle_0 + e^{-i(\omega_0/2 - N\bar{g})t}|1\rangle_0) \otimes_{i=1}^N |0\rangle_i$, up to an unimportant global phase factor. In particular, the state remains a product state at all times. If we measure $\sigma_x^{(0)}$ of the central spin, we find $\langle \sigma_x^{(0)}(t) \rangle = \cos((\omega_0 - 2N\bar{g})t)$, i.e. the oscillation frequency increases for large N proportional to N . As the quantum fluctuations of the central spin are independent of N , this implies a standard deviation in the estimation of the average coupling \bar{g} that scales as $1/N$, which can be confirmed by calculating the quantum Fisher information. Clearly, this is not an effect of entanglement, but simply of a phase accumulation. In this respect the approach is reminiscent to the sequential phase accumulation protocols, where the precision of a phase shift measurement is enhanced by sending the light several times through the same phase shifter (Higgins *et al.*, 2007). However, there the losses increase exponentially with the number of passes, and the sequential nature of the interaction leads to a bandwidth penalty that is absent for the simultaneous interaction described by H_{int} .

The fact that an interaction with N probes and a central quantum bus can lead to HL scaling of the sensitivity was first found in (Braun and Martin, 2011), where a more general model was studied. It will be noted that the total hamiltonian has exactly the same structure as for a decoherence model. However, in contrast to the standard scenario in decoherence, it is assumed here that at least part of the “environment” is accessible. The above example shows that the “environment” can be as simple as a single spin-1/2. One can also extend the model to include additional decoherence processes. In

(Braun and Martin, 2011) a phase-flip channel with rate Γ on all spins was considered. It was found that phase flips on any of the N probes has no effect, whereas phase flip of the central spin introduces a prefactor that decays exponentially as $\exp(-2\Gamma t)$ with time in $\langle \sigma_x^{(0)}(t) \rangle$. Since the prefactor is independent of N the power-law scaling of the sensitivity with N is unchanged, but it is clear that the prefactor matters and leads to a sensitivity that deteriorates exponentially with time.

The fact that the parameter estimated in the above example is the interaction strength between the bus and the N probes prevents a comparison with other schemes that do not use interactions. In (Fraïsse and Braun, 2015) a more comprehensive study of two different spin-systems was undertaken where also parameters describing the probes themselves and the bus were examined. Regimes of strong and weak interaction were analyzed, and different initial states considered. The two models, called ZZZZ and ZZXX models, are respectively given by the hamiltonians

$$H_1 = \frac{\omega_0}{2} \sigma_z^{(0)} + \frac{\omega_1}{2} \sum_{i=1}^N \sigma_z^{(i)} + \frac{g}{2} \sum_{i=1}^N \sigma_z^{(0)} \sigma_z^{(i)},$$

$$H_2 = \frac{\omega_0}{2} \sigma_z^{(0)} + \frac{\omega_1}{2} \sum_{i=1}^N \sigma_z^{(i)} + \frac{g}{2} \sum_{i=1}^N \sigma_x^{(0)} \sigma_x^{(i)}, \quad (59)$$

where $\hbar = 1$. The ZZZZ model is an exactly solvable dephasing model, the ZZXX can be analyzed numerically and with perturbation theory. The analysis is simplified when starting with a product state that is symmetric under exchange of the probes, in which case the probes can be considered as a single spin with total spin quantum number $j = N/2$. It was found that for ω_1 , HL scaling can be achieved in the ZZXX model when measuring the entire system, but not when only measuring the quantum bus, and not for the ZZZZ model. HL scaling for the uncertainty of ω_0 is not possible in either model. For the interaction strength g , HL scaling is found for a large set of initial states and range of interaction strengths when measuring the entire system, but only for a small set of initial states in the vicinity of the state considered in the simple example above, when measuring only the quantum bus. Interestingly, in the ZZZZ model HL scaling of the sensitivity for measuring g also arises with the probes in a thermal state at any finite temperature, as long as the quantum bus can be brought into an initially pure state. This is reminiscent of the “power of one-qubit” (Knill and Laflamme, 1998): With a set of qubits of which only a single one is initially in a pure state, a quantum enhancement is already possible in quantum computation, providing evidence for an important role of quantum discord (Datta *et al.*, 2005; Lanyon *et al.*, 2008) (see Sec. II.C). As detailed in Sec. II.C, the DQC1-scheme can provide better-than-SQL sensitivity as soon as the

ancillas have a finite purity. The control qubit plays the role of the quantum bus, and the dipole-dipole interaction between the Rydberg atoms implements the XX -interaction considered in (Fraïsse and Braun, 2015).

An important limitation of such schemes was proven in (Boixo *et al.*, 2007; Fraïsse and Braun, 2016), where it was shown very generally that with a hamiltonian extension to an ancilla system the sensitivity of a phase shift measurement cannot be improved beyond the best sensitivity achievable with the original system itself. Coherent averaging is nevertheless interesting, as it allows one to achieve without injecting entanglement better-than-SQL sensitivity for which the non-interacting phase shift measurement would need a highly entangled state.

The existence of HL scaling of sensitivities for product states, suggests that “coherent averaging” might even be possible classically. This question was investigated in (Braun and Popescu, 2014) in a purely classical model of harmonic oscillators, where N “probe” oscillators interact with a central oscillator. It was found that indeed for weak interaction strengths a regime of HL scaling of the sensitivity exists, even though the scaling crosses over to SQL scaling for sufficiently large N . Nevertheless, it was proposed that this could be useful for measuring very weak interactions, and notably improve measurements of the gravitational constant.

I. Quantum feedback schemes

In the context of having probes interact with additional (ancilla) systems, quantum feedback schemes should be mentioned. This is a whole field by itself (see e.g. (D’Alessandro, 2007; Serafini, 2012; Wiseman and Milburn, 2009) for reviews). Quantum feedback generalizes classical feedback loops to the quantum world: one tries to stabilize, or more generally dynamically control, a desired state of or an operation on a quantum system by obtaining information about its actual state or operation, and feeding back corrective actions into the controls of the system that bring it back to that desired state or operation if any deviation occurs. As a consequence, the field can be broadly classified according to two different categories: Firstly, the object to be controlled maybe a quantum state or an entire operation. And secondly, the type of information fed back can be classical or quantum. With classical information is meant information that is obtained from a measurement, and which is then typically processed on a classical computer and used to re-adjust the classical control knobs of the experiment. Such schemes are called “measurement-based feedback”. In contrast to this, “coherent feedback” schemes directly use quantum systems that are then manipulated unitarily and fed back to the system.

Measurement-based feedback schemes in metrology

are also known as “adaptive measurements” (see also Sec.V.D for adaptive measurements in the context of phase transitions). An adaptive scheme was proposed as early as 1988 by Nagaoka for mending the problem that the optimal POVM obtained in standard q-pet depends on the a priori unknown parameter (Nagaoka, 1988). One starts with a random estimate, uses its value to determine the corresponding optimal POVM, measures the POVM, updates the estimate, uses that new value to determine a new optimal POVM, and so on. The scheme was shown to be strongly consistent (meaning unbiased for the number of iterations going to infinity), and asymptotically efficient (i.e. saturates the QCRB in that limit) by Fujiwara (Fujiwara, 2006). It was experimentally implemented in (Okamoto *et al.*, 2012) as an adaptive quantum state estimation scheme for measuring polarization of a light beam, but is in principle a general purpose estimation scheme applicable to any quantum statistical model using identical copies of an unknown quantum state.

In the context of quantum optics, an adaptive homodyne scheme was proposed by Wiseman for measuring the phase of an optical mode, where the reference phase of the local oscillator $\Phi(t)$ is adjusted in real-time to $\Phi(t) \simeq \pi/2 + \varphi(t)$, where $\varphi(t)$ is the latest estimate of the phase carried by a continuous-wave, phase-squeezed light signal (Wiseman, 1995). This reference phase corresponds to the highest sensitivity in a homodyne scheme. Keeping the local oscillator phase through feedback close to this optimal operating point can beat non-adaptive heterodyning in single shot phase decoding, as experimentally demonstrated in (Armen *et al.*, 2002). Phase estimation using feedback was also studied in (Berry and Wiseman, 2000, 2002, 2006; Higgins *et al.*, 2007). In (Berry and Wiseman, 2002) it was investigated how well a stochastically varying, white noise-correlated phase can be estimated. The theoretical analysis showed that the variance of the phase estimation could be reduced by a factor $\sqrt{2}$ by a simple adaptive scheme compared to a non-adaptive heterodyne scheme, resulting in a value of $n^{-1/2}/\sqrt{2}$, where n is the number of photons per coherence time. With a squeezed beam and a more accurate feedback, the scaling can be improved to $n^{-2/3}$. The latter result was also found for a narrow-band squeezed beam (Berry and Wiseman, 2006, 2013). In (Yonezawa *et al.*, 2012) $15 \pm 4\%$ reduction of mean square error of the phase below the coherent-state limit was reported with this scheme in optical-phase tracking, i.e. in a case without any a priori information about the value of the signal phase. The broad support of the signal phase implies that there is an optimal amount of squeezing, and the sensitivity enhancement is directly given by the squeezing. The scheme can therefore be seen as a generalization of Caves’ idea of reducing the uncertainty with which a (fixed) phase shift in one arm of an interferometer can be measured

(Caves, 1981). Instead of having a fixed phase reference by the beam in the other arm of the interferometer, the feedback allows one to continuously adjust the phase of the reference in the homodyning scheme to the optimal operating point. In (Wheatley *et al.*, 2010) an “adaptive quantum smoothing method” was used experimentally for estimating a stochastically fluctuating phase on a coherent beam. “Smoothing” refers to the fact that estimates are obtained not only from data measured up to the time where one wants to estimate the phase, but also on later data. This implies, of course, that these smoothed estimates can only be calculated after a sufficient delay or at the end of the experiment, whereas feedback itself at time t can only use data from times $t' \leq t$ (or even $t' < t$ when considering finite propagation times). Theory predicts a reduction of the mean square error by a factor $2\sqrt{2}$ compared to the SQL (achievable by non-adaptive filtering, i.e. without feedback and using only previous data at any time), and an experimental reduction of about 2.24 ± 0.14 was achieved.

Quantum error correction (QEC) for quantum enhanced measurements, recently introduced in (Dür *et al.*, 2014; Kessler *et al.*, 2014), can also be seen in the context of quantum feedback schemes (Ahn *et al.*, 2002). QEC is one of the most important ingredients of quantum computing (Gottesman, 1996; Shor, 1995; Steane, 1996). The general idea, both for quantum computing and quantum enhanced measurements, is that one would like to apply recovery operations \mathcal{R} to a state that after encoding the desired information through an operation \mathcal{M} has been corrupted by a noise process \mathcal{E} , such that $\mathcal{R} \circ \mathcal{E} \circ \mathcal{M}(\rho) \propto \mathcal{M}\rho$. In (Kessler *et al.*, 2014) it was shown that this can be achieved for the sensing of a single qubit subject to phase noise if it is coupled to a pure ancilla bit. Syndrome measurements (i.e. measurements of collective observables which do not destroy the relevant phase information, a concept developed in QEC) of both qubits at a rate faster than the dephasing rate allow one to detect whether a phase flip has occurred and to correct it, extending in this way the coherence time available for Ramsey interferometry to much longer times and thus better maximum sensitivities. When using N spins-1/2, an ancilla is not necessary. The method then operates directly on an initially entangled state, such as the GHZ-state, and measures error-syndromes on pairs of spins. Thus, the idea is here not so much to avoid entanglement but rather to stabilize through rapid multi-spin error-syndrome measurements the correct imprinting of the information on the quantum state against unwanted decoherence processes.

V. THERMODYNAMICAL AND NON-EQUILIBRIUM STEADY STATES

In this section, we discuss precision parameter estimations with thermal states and non-equilibrium steady states of dissipative dynamics, which have the advantage to be stationary, and describe mesoscopic systems. Thermal probes are crucial for both fundamental issues and technological applications (Benedict, 1984; Childs, 2001; Giazotto *et al.*, 2006). Estimations with dissipative dynamics (Alipour *et al.*, 2014; Bellomo *et al.*, 2009, 2010a,b; Zhang and Sarovar, 2015) are also instances of process tomography (Baldwin *et al.*, 2014; Bendersky and Paz, 2013; Merkel *et al.*, 2013; Mohseni *et al.*, 2008) with partial prior knowledge. High sensitivity is found in correspondence of critical behaviors, e.g. separation between different states of matter or long-range correlations, providing precise probes, as is implicit in the identification of the QFI with the Bures metric. We lay emphasis on highly sensitive probes that do not need to be entangled, and in certain cases not even quantum.

A. Thermodynamical states and thermal phase transitions

Thermodynamic states at equilibrium are derived by the maximization of the information-theoretic Shannon’s entropy (Jaynes, 1957a,b), equivalent to the maximization of the number of microscopic configurations compatible with physical constraints. Given the probability distribution $\{p_j\}_j$ of a set of configurations $\{j\}_j$, the constraints are the normalization $\sum_j p_j = 1$ and the averages of certain quantities $\langle F^{(k)} \rangle = \sum_j p_j f_j^{(k)}$, $f_j^{(k)}$ being the values of the quantity $F^{(k)}$ corresponding to the j -th configuration. The solution of the maximization is the well-known Boltzmann-Gibbs distribution

$$p_j = \frac{e^{-\sum_k \theta_k f_j^{(k)}}}{Z}, \quad Z = \sum_j e^{-\sum_k \theta_k f_j^{(k)}}, \quad (60)$$

where Z is the partition function, and θ_k is the Lagrange multiplier corresponding to the quantity $F^{(k)}$ fixed on average.

This formalism is equally adequate for both classical and quantum thermodynamic systems. In the quantum case all the quantities $F^{(k)}$ are commuting operators, the configurations are labeled by the set of eigenvalues of these operators and possibly additional quantum numbers in the case of degeneracy, and the thermal state is the density matrix ρ diagonal in the common eigenbasis of the $F^{(k)}$ with eigenvalues p_j :

$$\rho = \sum_j p_j |j\rangle\langle j|, \quad F^{(k)}|j\rangle = f_j^{(k)}|j\rangle. \quad (61)$$

Lagrange multipliers are the thermodynamic parameters to be estimated. Since the density matrix in (61) depends on them only through its eigenvalues p_j , the quantum Fisher matrix $I = [I_{\theta_k, \theta_{k'}}]_{k, k'}$ coincides with the classical Fisher matrix of the probability distribution $\{p_j\}_j$. A straightforward computation shows

$$I_{\theta_k, \theta_{k'}} = \frac{\partial^2 \ln Z}{\partial \theta_k \partial \theta_{k'}} = \text{Cov}(F^{(k)}, F^{(k')}). \quad (62)$$

See also (Jiang, 2014) for the computation of the QFI with density matrices in exponential form. The diagonal element I_{θ_k, θ_k} is the largest inverse sensitivity for a single estimation of the parameter θ_k , while the Fisher matrix I bounds the inverse covariance matrix of the multiparameter estimation, see eq.(12). The Cramér-Rao bound hence reads

$$[\text{Cov}(\theta_{k, \text{est}}, \theta_{k', \text{est}})] \cdot [\text{Cov}(F^{(k)}, F^{(k')})] \geq \frac{1}{M}, \quad (63)$$

which is the uncertainty relation for conjugate variables in statistical mechanics (Davis and Gutiérrez, 2012; Gilmore, 1985).

The computation of the Fisher matrix (62), together with the Cramér-Rao bound (63), tells us that the best sensitivity of Lagrange multipliers $\{\theta_k\}_k$ is inversely proportional to squared thermal fluctuations, and thus susceptibilities (Reichl, 1998). For connections among metric of thermal states, Fisher information, and susceptibilities see (Brody and Rivier, 1995; Brody and Ritz, 2003; Crooks, 2007; Davis and Gutiérrez, 2012; Diósi *et al.*, 1984; Dolan, 1998; Janke *et al.*, 2003, 2002; Janyszek, 1986b, 1990; Janyszek and Mrugała, 1989; Mrugała, 1984; Nulton and Salamon, 1985; Prokopenko *et al.*, 2011; Ruppeiner, 1979, 1981, 1991, 1995; Salamon *et al.*, 1984; Weinhold, 1974) for classical systems and (Paunković and Vieira, 2008; Janyszek, 1986a; Janyszek and Mrugała, 1990; Marzolino and Braun, 2013, 2015; Quan and Cucchietti, 2009; You *et al.*, 2007; Zanardi *et al.*, 2007a, 2008) for quantum systems.

Due to the pairwise commutativity of the $F^{(k)}$, the estimations of the parameters $\{\theta_k\}_k$ are reduced to parameter estimations with the classical probability distribution $\{p_j\}_j$. Thus, the maximum likelihood estimator is asymptotically unbiased and optimal, in the sense of achieving the Cramér-Rao bound, in the limit of infinitely many measurements (Helstrom, 1976; Holevo, 1982). This estimator consists in measuring each $F^{(k)}$ with outcomes $\{f_j^{(k)}\}_{j=1, \dots, M}$, and in maximizing the average logarithmic likelihood $\ell = \frac{1}{M} \ln \prod_j p_j$ with respect to the parameters $\{\theta_k\}_k$.

Among the most common statistical ensembles for equilibrium systems, there are the canonical ensemble and the grandcanonical ensemble. The canonical ensemble describes systems that only exchange energy with

their surrounding: the only quantity $F_1 = H$ fixed on average is the Hamiltonian, the Lagrange multiplier is $\theta_1 = \beta = \frac{1}{k_B T}$ where k_B is the Boltzmann constant and T is the absolute temperature, and the Fisher information is proportional to the heat capacity $C_V = \left(\frac{\partial \langle H \rangle}{\partial T}\right)_V$,

$$I_{\beta, \beta} = \text{Var}(H) = k_B T^2 C_V. \quad (64)$$

The grandcanonical ensemble describes systems that exchange energy and particles with the surrounding: the quantities fixed on average are the Hamiltonian $F_1 = H$ and the particle number $F_2 = N$, with Lagrange multipliers being the inverse temperature $\theta_1 = \beta = \frac{1}{k_B T}$ and $\theta_2 = -\beta\mu$ where μ is the chemical potential. The Fisher information of temperature and chemical potential are linked to thermal fluctuations, i.e. heat capacity $C_V = \left(\frac{\partial \langle H \rangle}{\partial T}\right)_V$ and isothermal compressibility $\kappa_T = -\frac{1}{V} \left(\frac{\partial V}{\partial P}\right)_T$ respectively, where V is the volume and P is the pressure (Marzolino and Braun, 2013, 2015):

$$I_{\beta, \beta} = \text{Var}(\mu N - H) = \frac{\partial \mu \langle N \rangle}{\partial \beta} + k_B T^2 C_V \quad (65)$$

$$I_{\mu, \mu} = \beta^2 \text{Var}(N) = \frac{\langle N \rangle^2}{\beta V} \kappa_T. \quad (66)$$

Another parameter that can be estimated within this framework is the magnetic field. For certain classical magnetic or spin systems, the interaction with a magnetic field \mathbf{B} is $\mathbf{B} \cdot \mathbf{M}$, with \mathbf{M} being the total magnetization. This interaction term can represent a contribution to the Hamiltonian as well as additional “fixed-on-average quantities” \mathbf{B} with Lagrange multipliers $\beta \mathbf{M}$. The magnetic field is also linked to magnetic susceptibility $\chi = \frac{\partial \langle M \rangle}{\partial B}$ (where we consider for simplicity only a single component of \mathbf{B} and \mathbf{M} , $M \equiv M_z, \equiv Bz$):

$$I_{B, B} = \beta^2 \text{Var}(M) = \beta \chi, \quad (67)$$

This picture of the magnetic field as a Lagrange multiplier is valid also for coupling constants whenever the Hamiltonian is $H = \sum_j \lambda_j H_j$, where $\beta \lambda_j$ is the Lagrange multiplier of H_j . For general quantum systems, the non-commutativity of magnetization or other Hamiltonian contributions H_j with the rest of the Hamiltonian gives rise to QPTs that occur also at zero temperature without thermal fluctuations. The above considerations also apply to the so-called *generalized Gibbs ensembles*, i.e. with arbitrary fixed-on-average quantities $F^{(k)}$, where estimations of parameters θ_k are under experimental (Langen *et al.*, 2015) and theoretical (Foini *et al.*, 2016) study.

Thermal susceptibilities are typically extensive except in the presence of phase transitions. Thus, their connection with Fisher information suggests that thermal states at critical points (Baxter, 1982; Brody and Rivier,

1995; Brody and Ritz, 2003; Diósi *et al.*, 1984; Dolan, 1998; Janke *et al.*, 2003, 2002; Janyszek, 1990; Janyszek and Mrugała, 1989; Prokopenko *et al.*, 2011; Reichl, 1998; Ruppeiner, 1991, 1995) with divergent susceptibilities are probes for enhanced measurements. Thermal susceptibility divergences occur also in classical systems, proving precision measurements without entanglement.

1. Role of quantum statistics

We now discuss the estimation of Lagrange multipliers of quantum gases in the grandcanonical ensemble and the role of quantum statistics therein (Marzolino and Braun, 2013, 2015). Consider ideal gases in a homogeneous or harmonic trap. $I_{\beta,\beta}$ is always extensive in the average particle number $\langle N \rangle$. The corresponding relative error found in (Marzolino and Braun, 2013, 2015) for temperature estimation is still one order of magnitude smaller than the standard deviations obtained experimentally via density measurements of Bosons (Leanhardt *et al.*, 2003) and Fermions (Müller *et al.*, 2010; Sanner *et al.*, 2010).

Estimations of chemical potentials are more sensitive to quantum statistics than estimations of temperature, because the chemical potential is the conjugate Lagrange multiplier of the particle number which in turn reveals clear signatures of quantum statistics, such as bunching and antibunching. Effects of quantum statistics are more evident in quantum degenerate gases, i.e. at low temperatures.

In Fermion gases, $I_{\mu,\mu}$ is extensive but diverges at zero temperature. A change in the chemical potential corresponds to the addition or the subtraction of particles, thus achieving a state orthogonal to the previous one. This sudden state change makes the chemical potential estimation very sensitive. A generalization of the Cramér-Rao bound, called Hammersley-Chapman-Robbins-Kshirsagar bound that is suitable for non-differentiable statistical models (Tsuda and Matsumoto, 2005) must be applied. This may lead to superextensive $I_{\mu,\mu}$, depending on the degree of rotational symmetry breaking or confinement anisotropy and dimension (see appendix B in (Marzolino and Braun, 2013)).

Bose gases undergo a PT to a Bose-Einstein condensate (BEC) in three dimensions for homogeneous confinement, and in three or two dimensions in a harmonic trap. Approaching from above the critical temperature, or zero temperature with large density when there is no PT, $I_{\mu,\mu}$ is superextensive: for homogeneous and harmonic trap respectively

$$I_{\mu,\mu} \lesssim \begin{cases} \mathcal{O}(\beta^2 \langle N \rangle^{\frac{4}{3}}) & \text{in three dimensions} \\ \mathcal{O}(\beta^2 \frac{\langle N \rangle^2}{\log \langle N \rangle}) & \text{in two dimensions} \\ \mathcal{O}(\beta^2 \langle N \rangle^2) & \text{in one dimension} \end{cases}, \quad (68)$$

$$I_{\mu,\mu} \lesssim \begin{cases} \mathcal{O}(\beta^2 \langle N \rangle) & \text{in three dimensions} \\ \mathcal{O}(\beta^2 \langle N \rangle \log \langle N \rangle) & \text{in two dimensions} \\ \mathcal{O}(\beta^2 \frac{\langle N \rangle^2}{\log \langle N \rangle}) & \text{in one dimension} \end{cases}. \quad (69)$$

Below the critical temperature, the BEC phase depends on the anisotropy of the external potential. If the gas is much less confined along certain directions, the BEC is an effective low dimensional gas with excitations restricted to directions along the less confined dimensions. A hierarchy of condensations to subsequent lower-dimensional gases is possible. These BECs have been studied both at finite size and in the thermodynamic limit focusing on mathematical structures and general properties (Beau and Zagrebnov, 2010; van den Berg, 1983; van den Berg and Lewis, 1982; van den Berg *et al.*, 1986a,b; Casimir, 1968; van Druten and Ketterle, 1982; Girardeau, 1960, 1965; Ketterle and van Druten, 1982; Kreuger, 1968; Mullin, 1997; Mullin and Sakhel, 2012; Rehr, 1970; Zobay and Garraway, 2004), in connection with liquid helium in thin films (Douglass *et al.*, 1964; Goble and Trainor, 1965, 1966, 1967; Khorana and Douglass, 1965; Mills, 1964; Osborne, 1949), magnetic flux of superconducting rings (Sonin, 1969), and gravito-optical traps (Wallis, 1996). Experimental realizations of effective low dimensional BEC with trapped atoms have been reported in (van Amerongen, 2008; van Amerongen *et al.*, 2008; Armijo *et al.*, 2011; Bouchoule *et al.*, 2011; Esteve *et al.*, 2006; Görlitz *et al.*, 2001; Greiner *et al.*, 2001). $I_{\mu,\mu}$ in these BEC phases is superextensive and interpolates between the scaling above the critical temperature and $I_{\mu,\mu} = \mathcal{O}(\beta^2 \langle N \rangle^2)$ for standard BEC only consisting of the ground state. The advantage of BEC probes for precision estimations is that $I_{\mu,\mu}$ is superextensive in the entire BEC phase and not only at critical points as for susceptibilities of other thermal PT.

In a mean field model with interactions treated perturbatively, if the ideal system exhibits a superlinear scaling of the Fisher information, the interaction strength λ has to go to zero for $N \rightarrow \infty$ for the perturbation theory to remain valid. In this limit, the superlinear scaling disappears for any non-zero value of λ , but at finite N there are values of λ which do not destroy the sub-shot-noise.

Moreover, superextensive QFI results in one dimension at fixed volume L_x and small contact interaction coupling

$$\begin{aligned}
I_{\mu,\mu} \simeq & \frac{\beta^2 \lambda_T^2 \langle N \rangle^3}{2\pi L_x^2} + \beta^2 \langle N \rangle - \frac{\beta^2 \lambda_T^4 \langle N \rangle^4}{8\pi^2 L_x^4} \left(1 - e^{-\frac{4\pi}{\lambda_T^2 \rho^2} \langle N \rangle} \right) \\
& + c \left(\frac{3\beta^3 \lambda_T^6 \langle N \rangle^7}{16\pi^3 L_x^7} \left(1 - e^{-\frac{4\pi}{\lambda_T^2 \rho^2} \langle N \rangle} \right) \right. \\
& \left. - \frac{\beta^3 \lambda_T^4 \langle N \rangle^6}{4\pi^2 L_x^4} \left(2 + e^{-\frac{4\pi}{\lambda_T^2 \rho^2} \langle N \rangle} \right) \right), \quad (70)
\end{aligned}$$

where λ_T is the thermal wavelength, in agreement with experimental measurements on atom chips using ^{87}Rb atoms (Armijo *et al.*, 2011). Superextensive grand-canonical fluctuations of particle number, and thus superextensive $I_{\mu,\mu}$, have been observed in photon BEC (Schmitt *et al.*, 2014) which can be realized at room temperature (Klaers, 2014).

2. Interferometric thermometry

We now discuss a protocol for precision thermometry proposed in (Stace, 2010), using a Mach-Zehnder interferometer coupled to an ideal gas consisting of N two-level atoms in the canonical ensemble. The gas Hamiltonian is $H_0 = \sum_{j=1}^N \epsilon |\epsilon\rangle_j \langle \epsilon|$, where $|\epsilon\rangle_j$ is the j -th particle excited state with single-particle energy ϵ , while the single-particle lowest energy is zero. We skip the label j at the bra-vector in the projector for brevity.

The interferometer is injected with K two-level atoms that interact with the gas in one arm of the interferometer with the interaction Hamiltonian $H_I = \alpha \sum_{j=1}^K \sum_{l=1}^N |\epsilon\rangle_j \langle \epsilon| \otimes |\epsilon\rangle_l \langle \epsilon|$, where the index j labels the atoms in the interferometer and l refers to the atoms in the gas. In order the interaction not to sensitively perturb the gas, the interferometer should be much smaller than the gas, thus $K \ll N$. Each atom in the interferometers acquires a relative phase $\phi = \alpha m \tau$ between the arms, where τ is the interaction time and m is the number of excited atoms in the gas whose expectation $\langle m \rangle = N/(1 + e^{\beta\epsilon})$ depends on the temperature.

The inverse temperature can be estimated from the interferometric phase measurement with the sensitivity

$$\sigma(\beta) = \frac{\delta\phi}{\left| \alpha \tau \frac{d\langle m \rangle}{d\beta} \right|} = \frac{(1 + e^{\beta\epsilon})^2}{\epsilon e^{\beta\epsilon} N} \cdot \frac{\delta\phi}{\alpha \tau}, \quad (71)$$

resulting from error propagation, where $\delta\phi$ is the best sensitivity of the phase estimation according to the QCRB. The scaling with the number K of probes comes from $\delta\phi$. For distinguishable atoms, separable states imply shot-noise $\delta\phi = 1/K^{1/2}$, while sub-shot-noise can be achieved with separable states of identical atoms as discussed in section III, leading to $\delta\beta \propto 1/K$ for ideal non-interacting bosons, or $\delta\beta \propto 1/K^{p+1/2}$ for ideal fermions with a dispersion relation of the probe atoms

as discussed after (45). Note also that the scaling of the sensitivity (71) with respect to the particle number N in the gas looks like a Heisenberg scaling irrespectively of the probe state, and it is not possible to achieve this scaling by direct measurement of the gas (i.e. without any probe) because the Fisher information is extensive except at critical points. What is conventionally considered is however the scaling with the number of probes, in this case K , which can be controlled. Optimal thermometry with a single quantum probe (and hence by definition without entanglement) was discussed in (Correa *et al.*, 2015). It was found that the QCRB for T of a thermalizing probe reproduces the well-known relation of temperature fluctuations to specific heat $C_V(T)$, $(T/\text{Var}(T_{\text{est}}))^2 \leq C_V(T)$ (with Boltzmann constant $k_B = 1$, see also (Jahnke *et al.*, 2011)). The level structure of the probe was then optimized to obtain maximum heat capacity and it was found that the probe should have only two different energy levels, with the highest one maximally degenerate and a non-trivial dependence of the optimal energy gap on temperature (see also (Reeb and Wolf, 2015)). The sensitivity of the probe increases with the number of levels, but the role of the quantumness of the initial state of the probe on the sensitivity of thermometry is not fully understood yet. Interferometric thermometry with a single probe was realized experimentally in an NMR setup in (Raitz *et al.*, 2015) and the role of quantum coherence emphasized. Experimental simulations in quantum optical setups and investigations of the role of quantum coherence were reported in (Mancino *et al.*, 2016; Tham *et al.*, 2016). The theoretical study in (Jevtic *et al.*, 2015) examined thermometry with two qubits. Numerical evidence suggested that while initial quantum coherences can improve the sensitivity, the optimal initial state is not maximally entangled.

B. Thermodynamical states and quantum phase transitions

QPTs are sudden changes of the ground state for varying Hamiltonian parameters. If the Hamiltonian $H(\boldsymbol{\theta}) = \sum_{j \geq 0} E_j(\boldsymbol{\theta}) |E_j(\boldsymbol{\theta})\rangle \langle E_j(\boldsymbol{\theta})|$, with $E_j \leq E_{j+1}$, has a unique pure ground state $|E_0(\boldsymbol{\theta})\rangle$, the Fisher matrix (Campos Venuti and Zanardi, 2007; Gu and Lin, 2009; You *et al.*, 2007; Zanardi *et al.*, 2007c) is

$$I_{\theta_k, \theta_{k'}} = 4 \text{Re} \sum_{j>0} \frac{\langle E_0 | \partial_{\theta_k} H | E_j \rangle \langle E_j | \partial_{\theta_{k'}} H | E_0 \rangle}{(E_j - E_0)^2} \quad (72)$$

which follows from the differentiation of the eigenvalue equation $H|E_j\rangle = E_j|E_j\rangle$ or from the standard time-independent perturbation theory with respect to small variations $d\boldsymbol{\theta}$. QFI is also expressed in terms of the imaginary time correlation function or dynamical response function (Campos Venuti and Zanardi,

2007; Gu and Lin, 2009; You and He, 2015; You *et al.*, 2007), and has been used to count avoided crossings (Wimberger, 2016).

Equation (72) tells us that the Fisher matrix can diverge only for divergent Hamiltonian derivatives or for gapless systems $E_1 - E_0 \rightarrow 0$ in the thermodynamic limit. Thus, the divergence or the superextensivity of the QFI reveals a QPT but the converse does not hold: see (Gu, 2010) for a review. Finite size scaling of the QFI (Campos Venuti and Zanardi, 2007; Gu and Lin, 2009; Zanardi *et al.*, 2008) can be derived using finite size scaling at criticality (Brankov *et al.*, 2000; Continentino, 2001).

Superextensive QFI was observed at low order symmetry breaking QPT, topological QPT, and gapless phases, but this criterion may fail at high order symmetry breaking QPT (Tzeng *et al.*, 2008) and Berezinskii-Kosterlitz-Thouless (BKT) QPT (Chen *et al.*, 2008; Sun *et al.*, 2015). We now review quantum critical systems at zero temperature exhibiting superextensive QFI of Hamiltonian parameters without entanglement. We adapt and unify the notation of existing literature.

1. Quasi-free Fermion models

We start with quasi-free Fermion models (Cozzini *et al.*, 2007), i.e. quadratic Hamiltonians in the creation c_j^\dagger and annihilation c_j operators of L Fermionic modes:

$$H_{\text{quasi-free}} = \sum_{j,l=1}^L \left(c_j^\dagger A_{j,l} c_l + \frac{1}{2} \left(c_j^\dagger B_{j,l} c_l^\dagger + \text{h.c.} \right) \right), \quad (73)$$

where L is a measure of the system volume. The Hamiltonian (73) is transformed by a Bogoliubov rotation into an ideal Fermion Hamiltonian in terms of new modes, hereafter called “eigenmodes”. The eigenstates of (73), and thus thermal probes including the ground state, are not entangled in the eigenmodes even though they are typically entangled in the modes $\{c_j, c_j^\dagger\}_j$, providing precision estimations without entanglement. The relativity of entanglement with respect to the basis of modes and complementary pictures of estimation protocols is reminiscent of the case of interferometry with identical particles discussed in Section III.

Translationally invariant Hamiltonians (73) with periodic boundary conditions and tunneling in the modes $\{c_j, c_j^\dagger\}_j$ of range r have

$$\begin{aligned} A_{j,l} &= (J - \mu)\delta_{j,l} - J\theta(r - |j - l|) - J\theta(|j - l| - L + r), \\ B_{j,l} &= J\gamma \text{sign}(j - l)(\theta(r - |j - l|) - \theta(|j - l| - L + r)), \end{aligned} \quad (74)$$

where $J > 0$, $\text{sign}(0) = 0$, and $\theta(\cdot)$ is the unit step function with $\theta(0) = 1$. Here J corresponds to a tunneling energy between different sites, and $J\gamma$ to an effective interaction; μ multiplies the total particle number and hence corresponds to a chemical potential. This Hamiltonian can be analytically diagonalized (Cozzini *et al.*, 2007). For large L and a fully connected system, $r = \lfloor \frac{L}{2} \rfloor$ with periodic boundary conditions, the Fisher matrix with respect to (μ, γ) is

$$I_{\mu,\mu} = \gamma^2 S, \quad I_{\gamma,\gamma} = (\mu - J)^2 S, \quad I_{\mu,\gamma} = -(\mu - J)\gamma S, \quad (75)$$

with

$$S \simeq \begin{cases} \frac{L}{2} \left| \frac{J}{(\mu - J)\gamma} \right| \frac{1}{(|\mu - J| + |J\gamma|)^2} & \text{if } (\mu \neq J, \gamma \neq 0) \\ \frac{L^2}{3J^2\gamma^4} & \text{if } (\mu = J, \gamma \neq 0) \\ \frac{L^2 J^2}{(\mu - J)^4} & \text{if } (\mu \neq J, \gamma = 0) \end{cases} \quad (76)$$

The lines $(\mu = J, \gamma \neq 0)$ and $(\mu \neq J, \gamma = 0)$ reveal second-order QPTs, as well as superextensivity of the Fisher matrix in the volume L . At the critical line $\gamma = 0$, the ground state is unentangled not only in the eigenmodes but also in the original modes $\{c_j, c_j^\dagger\}_j$.

The fully connected translationally invariant Hamiltonian with open boundary conditions reads

$$A_{j,l} = (\mu - J)\delta_{j,l} + J, \quad B_{j,l} = J\gamma \text{sign}(l - j), \quad (77)$$

with second-order QPTs at lines $\mu = J$ and $\gamma = 0$, and superextensive Fisher matrix as for periodic boundary conditions with different prefactors (Cozzini *et al.*, 2007; Zanardi *et al.*, 2007b). For instance, at $(\mu > J, \gamma = 0)$

$$I_{\mu,\mu} = 0, \quad I_{\gamma,\gamma} = \frac{L^2 J^2}{3(\mu - J)^2}, \quad I_{\mu,\gamma} = 0. \quad (78)$$

Another interesting case is the Hamiltonian with nearest neighbor tunneling in the modes $\{c_j, c_j^\dagger\}_j$, periodic boundary conditions, and $J > 0$:

$$\begin{aligned} A_{j,l} &= (J - \mu)\delta_{j,l} - J\theta(1 - |j - l|), \\ B_{j,l} &= J\gamma \text{sign}(l - j)\theta(1 - |j - l|), \end{aligned} \quad (79)$$

Such Hamiltonian (Zanardi *et al.*, 2008) is also equal to

$$H_{\text{quasi-spin}} = \sum_{n=1}^{\lfloor \frac{L}{2} \rfloor} (-\epsilon_n \sigma_n^z + \Delta_n \sigma_n^y), \quad (80)$$

with $\epsilon_n = -J \cos(\frac{2\pi n}{L}) - \frac{\mu}{2}$, $\Delta_n = -J\gamma \sin(\frac{2\pi n}{L})$, and $\{\sigma_n^{y,z}\}_n$ are Pauli matrices on n orthogonal \mathbb{C}^2 subspaces. Eq. (80) provides an alternative representation of the Hamiltonian in terms of non-interacting quasi-spins, and its eigenstates are separable with respect to both the eigenmodes and the quasi-spins. The same QFI scaling holds in both representations as theoretical independent models.

QFI is linear in L away from the critical points, but superextensive at criticality in the leading order for large L (Zanardi and Paunković, 2006; Zanardi *et al.*, 2008):

$$I_{J,J}(|\mu| = 2J, \gamma) = \mathcal{O}\left(\frac{L^2}{J^2\gamma^2}\right), \quad (81)$$

$$I_{\gamma,\gamma}(|\mu| \leq 2J, \gamma = 0) = \mathcal{O}(L^2). \quad (82)$$

The origin of the superextensivity of $I_{J,J}$ stems from the fact that the symmetric logarithmic derivative, and thus the optimal estimation of J , is close to a single particle operator in the Fermion representation away from the critical point, but is a genuine multi-particle operator close to the critical point. In the canonical ensemble, at $|\frac{\mu}{2} - J| \lesssim \frac{1}{\beta}$, the QFI has a divergence around zero temperature, $I_{J,J} = \mathcal{O}(\frac{L\beta}{J\gamma})$. The superextensivity of $I_{J,J}$, together with the divergence of the derivative of the geometric phase, is also a universal feature depending only on the slope for the closing of the energy gap between one and zero Fermion occupations of certain eigenmodes (Cheng *et al.*, 2016).

Superextensive QFI is observed also for tight-binding electrons on the triangular lattice with magnetic flux $\frac{\phi}{2}$ within each triangle, hopping constants t_a (t_b) at edge along the x (y) direction and t_c at the third edge (Gong and Tong, 2008). Assuming zero momentum in the y direction (Ino and Kohmoto, 2006), the Hamiltonian can be transformed into (73), with $B_{j,l} = 0$ and

$$A_{j,l} = -2t_b \cos(2\pi\phi j) \delta_{l,j} - \left(t_a + t_c e^{-2\pi i\phi(j-\frac{1}{2})}\right) \delta_{l,j-1} - \left(t_a + t_c e^{2\pi i\phi(j+\frac{1}{2})}\right) \delta_{l,j+1}, \quad (83)$$

with j and l labeling the sites in the x direction. Consider $L = F_m$, the m -th Fibonacci number, and $\phi = \frac{F_{m-1}}{F_m} \xrightarrow{m \rightarrow \infty} \frac{\sqrt{5}-1}{2}$, the inverse of the golden ratio in the thermodynamic limit. At the critical line $t_c = t_a$, numerical computations result in $I_{t_c/t_a} = \mathcal{O}(L^{4.9371})$ if $m = 3L + 1$, otherwise $I_{t_c/t_a} = \mathcal{O}(L^{2.0})$. I_{t_b/t_a} has the same size dependence at the critical line $2t_b = t_a$. A more complicated quasi-free Fermion Hamiltonian describes a superconductor with a magnetic impurity (Paunković *et al.*, 2008). A sudden increase of the Bures distance between two reduced ground states of few modes around the impurity at very close exchange interaction with the impurity was numerically observed, but

it is unclear whether the QFI with respect to exchange interaction is superextensive.

2. Hubbard models

Another class of Fermionic systems is described by L -mode Hubbard Hamiltonians

$$H_{\text{HM}} = -\sum_{\substack{j=1,\dots,L \\ \sigma=\uparrow,\downarrow}} t_\sigma \left(c_{j,\sigma}^\dagger c_{j+1,\sigma} + \text{h.c.} \right) + U \sum_{j=1}^L n_{j,\uparrow} n_{j,\downarrow} - \mu \sum_{j,\sigma} n_{j,\sigma}, \quad (84)$$

with $n_{j,\sigma} = c_{j,\sigma}^\dagger c_{j,\sigma}$, t_σ the hopping constants, U the interaction strength, and μ the external potential. When $t_\uparrow = t_\downarrow$, the system undergoes a Berezinskii-Kosterlitz-Thouless (BKT) QPT at $U = 0$ and half filling $n = \frac{1}{L} \sum_{j,\sigma} \langle n_{j,\sigma} \rangle = 1$. I_U is extensive, but I_U/L diverges as $1/n$ for $n \rightarrow 0$ and $U = 0$, and as $1/U^4$ around the BKT critical point only if the system size is much larger than the correlation length (Campos Venuti *et al.*, 2008). At zero interaction (84) is a free Fermion Hamiltonian, and thus the ground state is not entangled in the eigenmodes.

In the large U limit and at $n = \frac{2}{3}$, a QPT occurs with control parameter t_\downarrow/t_\uparrow . QFI $I_{t_\downarrow/t_\uparrow}$ at critical points is superextensive $\mathcal{O}(L^\alpha)$, where the exponent was numerically computed $\alpha \simeq 5.3$ (Gu *et al.*, 2008). In the large U limit, the eigenstates of the Hamiltonian are perturbations of those of the interaction term, that are Fock states, thus with vanishingly small entanglement with respect to the modes $\{c_{j,\sigma}, c_{j,\sigma}^\dagger\}_{j,\sigma}$.

3. Spin-1/2 systems

Consider now systems of N spins-1/2 and Hamiltonian

$$H_{\text{spin}} = -J \sum_{j=1}^{N-1} \left(\frac{1+\gamma}{2} \sigma_j^x \sigma_{j+1}^x + \frac{1-\gamma}{2} \sigma_j^y \sigma_{j+1}^y + \Delta \sigma_j^z \sigma_{j+1}^z \right) + d \sum_{j=1}^{N-1} (\sigma_j \wedge \sigma_{j+1})^z - \sum_{j=1}^N (h - gj) \sigma_j^z, \quad (85)$$

anisotropies γ and Δ , Dzyaloshinskii-Moriya coupling d , magnetic field with uniform value h and gradient g , and $\sigma_j = (\sigma_j^x, \sigma_j^y, \sigma_j^z)$.

For the moment, we focus on the XY model with transverse field, i.e. $\Delta = d = g = 0$, with $J > 0$ and periodic boundary conditions. This Hamiltonian can be transformed into (73) with (79), $\mu = 2h$, and $L = N$ using the Jordan-Wigner transformation (Giamarchi, 2003). Thus, the XY Hamiltonian provides an alternative physical setting to implement precision metrology without entanglement in the Fermionic eigenmodes.

At zero temperature, $I_{h,h} = (\frac{J^2}{h^2})I_{J,J}$, since the ground state depends on h and J only via the ratio $\frac{h}{J}$. The Fisher matrix is extensive with divergent prefactors around critical regions (Cheng *et al.*, 2016; Zanardi *et al.*, 2007c), and is superextensive at criticality as in (81,82).

In the Ising model, i.e. $\gamma = 1$, corrections to the QFI around the critical points $|h| = J$ (Chen *et al.*, 2008; You and He, 2015; Zhou and Barjaktarević, 2008; Zhou *et al.*, 2008) and at small temperature (Invernizzi *et al.*, 2008; Zanardi *et al.*, 2007a) are also linear in N with divergent prefactors. $I_{h,h}$ was also explicitly computed in (Damski, 2013; Damski and Rams, 2014; You and He, 2015). A good estimator of the parameter J is the value inferred by measurements of the total magnetization, at least at small system size (Invernizzi *et al.*, 2008).

The divergence of $I_{h,h}/N$ was numerically observed also in the XX model, i.e. $\gamma = 0$, approaching the critical field $|h| = J$, with the reduced state of spin blocks (Sacramento *et al.*, 2011). This implies precision estimations of the magnetic field looking only at a part of the system.

Divergent $I_{h,h}/N$ at critical field $|h| = J$ was numerically computed also when an alternating magnetic field $\sum_{j=1}^N (-)^{j+1} \delta \sigma_j^z$ is added to the XY Hamiltonian (You and He, 2015). The corresponding quasi-free Fermion Hamiltonian has the additional term $\sum_{j=1}^N (-)^{j+1} \delta c_j^\dagger c_j$.

The XY and the Ising models are also interesting because the superextensivity of the Fisher matrix $I = [I_{\theta,\theta'}]_{\theta,\theta'=h,\gamma}$ is robust against disorder, i.e. with Hamiltonian parameters being Gaussian random variables (Garnerone *et al.*, 2009b). In the presence of disorder, QPTs are broadened to Griffiths phases (Fisher, 1992, 1995; Griffiths, 1969; Sachdev, 1999). Thus, the exponent of N in average $I_{h,h}$ and $I_{\gamma,\gamma}$ is slightly reduced but the superextensivity is broadened in the parameter range. Superextensivity of $I_{h,h}$ at the critical Ising point $|h| = J$ becomes a broad peak, and the superextensive $I_{\gamma,\gamma}$ at the $\gamma = 0$ critical line is split into two symmetric broad peaks around $\gamma = 0$, where at $\gamma = 0$ a local minimum with extensive $I_{\gamma,\gamma}$ arises. Furthermore, the scaling of $I_{h,h}$ of the XY model at the critical point $|h| = J$ is preserved when a periodic time-oscillating transverse magnetic field is considered, and the state is the Floquet time-evolution of the ground state at $t = 0$ up to times that scale linearly with the system size for small time-dependent driving (Lorenzo *et al.*, 2016). These robustness features are suitable for practical implementations of precision measurements.

Another interesting system is the XXZ model, i.e. $\gamma = 0$ and $J < 0$, whose low-energy spectrum for $|\Delta| < \frac{1}{2}$ is equivalent to a quasi-free Boson Hamiltonian known as Luttinger liquid (Giamarchi, 2003). Thus, superextensive ground state QFI implies precision measurements without entanglement in the Boson eigenmodes. Given

the Luttinger liquid parameter $K = \frac{\pi}{2 \arccos(-\frac{\pi}{\sqrt{1+2d/J}})}$, superextensive QFI I_d is observed with periodic boundary conditions at $d = h = g = 0$,

$$I_d = \begin{cases} \mathcal{O}\left(\frac{N \ln N}{J^2}\right) & \text{if } \Delta = \frac{1+\sqrt{5}}{8} \\ \mathcal{O}\left(\frac{N^{6-8K}}{J^2}\right) & \text{if } \frac{1+\sqrt{5}}{8} < \Delta < \frac{1}{2}, \\ \mathcal{O}\left(\frac{N}{J \ln N}\right)^2 & \text{if } \Delta = \frac{1}{2} \end{cases}, \quad (86)$$

while, with open boundary conditions, $I_d = \mathcal{O}\left(\frac{J^2 K N^2}{(J^2 + 4d^2)^2}\right)$ at $h = g = 0$ and $I_g = \mathcal{O}\left(\frac{K N^4}{J^2 u^2}\right)$, with $u = \frac{\pi \sqrt{1-4\Delta^2}}{2 \arccos(2\Delta)}$, at $d = h = g = 0$ (Greschner *et al.*, 2013). Here, QFI is superextensive mainly because the system is gapless, while the QPT only affects subleading in N , even though superextensive, orders of QFI. I_d is also equal to the QFI I_ϕ with respect to a twist phase ϕ of spin operators $\sigma_j^+ \sigma_{j+1}^- \rightarrow \sigma_j^+ \sigma_{j+1}^- e^{i\phi}$ (Thesberg and Sørensen, 2011). Via the Jordan-Wigner transformation, the $\Delta = \gamma = 0$ case in (85) is equivalent to a quasi-free Fermion Hamiltonian with matrix elements given by (79) with $\gamma = \mu = 0$ and with the Dzyaloshinskii-Moriya and the gradient field terms $-\frac{id}{2} \sum_{j=1}^N c_j^\dagger c_{j+1} + \text{h.c.}$ and $\sum_{j=1}^N g j c_j^\dagger c_j$, respectively. I_d and I_g have the same scaling in a generalization of the latter model with spin-1/2 Fermions (Greschner *et al.*, 2013), providing another example of precision measurements without Fermion entanglement.

Another model, mapped to a quasi-free Fermionic Hamiltonian and exhibiting superextensive QFI, is the quantum compass chain with periodic boundary conditions (Motamedifar *et al.*, 2013)

$$H_{\text{QCC}} = \sum_{j=1}^{\frac{N}{2}} \left(\sum_{\alpha=x,y} J_\alpha \sigma_{2j-1}^\alpha \sigma_{2j}^\alpha + J_z \sigma_{2j}^z \sigma_{2j+1}^z \right) - h \sum_{j=1}^N \sigma_j^y. \quad (87)$$

QPTs occur at critical fields $h_{1,2} = \sqrt{J_x(J_y \pm J_z)}/2$ when these values are real. Numerical computations show that $I_h = \mathcal{O}(N^\alpha)$ is superextensive with

$$\alpha \simeq \begin{cases} 1.80 \pm 0.02 & \text{for } \frac{J_x}{J_z} > 0, \frac{J_y}{J_z} > 1 \text{ and } h = h_1 \\ 1.98 \pm 0.02 & \text{for } \frac{J_x}{J_z} > 0, \frac{J_y}{J_z} > 1 \text{ and } h = h_2 \\ 1.94 \pm 0.02 & \text{for } \frac{J_x}{J_z} < 0, \frac{J_y}{J_z} < 1 \text{ and } h = h_2 \\ 2.02 \pm 0.02 & \text{for } \frac{J_x}{J_z} > 0, \frac{J_y}{J_z} < 1 \text{ and } h = h_1 \end{cases}. \quad (88)$$

4. Topological quantum phase transitions

QFI is superextensive also in topological QPT which have non-local order parameters (Zeng *et al.*, 2015). Mosaic models defined on two-dimensional lattices with trivalent vertices, i.e. each vertex is the border among

three polygons, and with three-body interaction were numerically studied. The N -particle Hamiltonian is

$$H_{\text{mosaic}} = - \sum_{\substack{\alpha=x,y,z \\ (j,l) \in S(\alpha)}} J_{j,l}^{\alpha} \sigma_j^{\alpha} \sigma_l^{\alpha} - K \sum_{j,l,k} \sigma_j^x \sigma_l^y \sigma_k^z, \quad (89)$$

where $S(\alpha)$ is the set of edges in the $\alpha \in \{x, y, z\}$ direction. H_{mosaic} can be mapped onto a free Majorana Fermion Hamiltonian, thus without entanglement in Fermion eigenmodes. When the edge numbers of the three polygons are $(4, 8, 8)$, $J_{j,l}^{x,y} = J$, $J_{j,l}^{z} = J_z$, and $K = 0$, the QFI is $I_{J_z} = \mathcal{O}(N^{1.07615 \pm 0.00005})$ at the critical point $J_z = \sqrt{2}J > 0$ (Garnerone *et al.*, 2009a). When the edge numbers are $(3, 12, 12)$, $J_{j,l}^{x,y,z} = J$ for edges within triangular elementary subcells (Yao and Kivelson, 2007), $J_{j,l}^{x,y,z} = J'$ for other links, and $K = 0$, the QFI is $I_{J'} = \mathcal{O}(N^{1.078 \pm 0.005})$ at the critical point $J' = \sqrt{3}J > 0$ (Garnerone *et al.*, 2009a). H_{mosaic} on the honeycomb lattice with $J_{j,l}^{\alpha} = J_{\alpha}$ has topological QPT at the boundaries $|J_x| = |J_y| + |J_z|$, $|J_y| = |J_z| + |J_x|$, and $|J_z| = |J_x| + |J_y|$ with superextensive QFI, e.g. $I_{J_x} = \mathcal{O}(N \ln N)$ for $J_y = J_z = \frac{J_x}{2}$ and $K = 0$ (Zhao and Zhou, 2009), and $I_{J_z} = \mathcal{O}(N^{1.08675 \pm 0.00005})$ for $J_z = \frac{J}{2}$, $J_x = J_y = \frac{J-J_z}{2}$ and $K = \frac{1}{15}$ (Garnerone *et al.*, 2009a).

Topological QPT are particularly relevant because of superextensive QFI in gapless phases and not only at critical points. In the honeycomb lattice with $K = 0$, $J_x + J_y + J_z = J$ and $J_x = J_y$, the QFI is $I_{J_z} = \mathcal{O}(N^{1.2535 \pm 0.00005})$ at the critical point $J_z = \frac{J}{2}$ (Yang *et al.*, 2008), and $I_{J_z} = \mathcal{O}(N \ln N)$ in the gapless phase $J_z < \frac{J}{2}$ (Gu and Lin, 2009). Superextensivity in the gapless phase is due to the algebraic decay of the correlation function of the z -edge, in contrast to exponential decay in the gapped phase $J_z > \frac{J}{2}$ leading to extensive QFI.

C. Non-equilibrium steady states

Consider now parameter estimation with probes in non-equilibrium steady states (NESS) of spin-1/2 chains with boundary noise (Marzolino and Prosen, 2016a; Prosen, 2015; Zunkovic and Prosen, 2010) described by the Markovian master equation (Benatti and Floreanini, 2005; Breuer and Petruccione, 2002)

$$\frac{\partial \rho_t}{\partial t} = -i[H_{\text{XYZ}}, \rho_t] + \sum_{\substack{\alpha=1,2 \\ j=1,N}} \left(\mathcal{L}_{\alpha,j} \rho_t \mathcal{L}_{\alpha,j}^{\dagger} - \frac{1}{2} \left\{ \mathcal{L}_{\alpha,j}^{\dagger} \mathcal{L}_{\alpha,j}, \rho_t \right\} \right), \quad (90)$$

with XYZ Hamiltonian, i.e. (85) with $d = g = 0$, and Lindblad operators $\mathcal{L}_{1(2),j} = \sqrt{(1 \pm \mu)/2} \sigma_j^{\pm}$.

Superextensive QFI of the NESS ρ_{∞} is observed at non-equilibrium PT and in phases with long-range correlations. In the XY model, $\Delta = 0$, superextensive QFI was

computed at the critical lines $h = 0$ ($I_{h,h} = \mathcal{O}(N^6)$) and $\gamma = 0$ with $|h| < J|1 - \gamma^2|$ ($I_{\gamma,\gamma} = \mathcal{O}(N^2)$), at the critical points $|h| = J|1 - \gamma^2|$ ($[I_{\theta,\theta'}]_{\theta,\theta'=h,\gamma} = \mathcal{O}(N^6)$), and in the phase with long-range correlations $|h| < J|1 - \gamma^2|$ ($[I_{\theta,\theta'}]_{\theta,\theta'=h,\gamma} = \mathcal{O}(N^3)$) (Banchi *et al.*, 2014).

In the XXZ model $\gamma = 0$, QFI I_{Δ} is superextensive in the limit of small $\frac{\lambda}{J}$ and for $|\Delta| \leq \frac{1}{2}$ at irrational $\frac{\arccos \Delta}{\pi}$ (Marzolino and Prosen, 2014, 2016b, 2017). If $\frac{\arccos \Delta}{\pi}$ is rational $I_{\Delta} = \frac{\lambda^2 \mu^2}{J^2} (\tilde{\xi} N^2 + \xi N)$ where ξ and $\tilde{\xi}$ are constants in N and for $\frac{\lambda}{J} < \frac{1}{\sqrt{N}}$, thus the QFI is not superextensive. Nevertheless, after inserting the value of Δ in $\frac{\arccos \Delta}{\pi}$ and reducing the fraction to lowest term $\frac{\arccos \Delta}{\pi} = \frac{q}{p}$ with coprime integers q, p , one realizes that the coefficient ξ is unbounded when the denominator p grows. Therefore, I_{Δ} , as a function of Δ , exhibits a fractal-like structure with a different size scaling at irrational $\frac{\arccos \Delta}{\pi}$. Moreover, the limit of $\frac{\arccos \Delta}{\pi}$ approaching irrational numbers from rationals and the thermodynamic limit do not commute. If the thermodynamic limit is first performed then $I_{\Delta} = \mathcal{O}\left(\frac{\lambda^2 \mu^2}{J^2} N^5\right)$ with $\frac{\lambda}{J} < \frac{1}{\sqrt{N}}$. If the thermodynamic limit is postponed after the limit to irrational $\frac{\arccos \Delta}{\pi}$ then QFI is still superextensive with a power law depending on the value of Δ , e.g. $I_{\Delta} = \mathcal{O}\left(\frac{\lambda^2 \mu^2}{J^2} N^{2.3}\right)$ with $\frac{\lambda}{J} < \frac{1}{\sqrt{N}}$ and $\frac{\arccos \Delta}{\pi}$ being the golden ratio. When $\Delta = \frac{1}{2}$, $I_{\Delta} = \mathcal{O}\left(\frac{\lambda^2 \mu^2}{J^2} N^4\right)$ with $\frac{\lambda}{J} < \frac{1}{\sqrt{N}}$ in both cases.

In addition, the QFI of the reduced state of a single spin at position k scales superextensively also for arbitrary dissipation strength λ but only at the critical points $|\Delta| = 1$, and with a power law depending on the position k of the spin (Marzolino and Prosen, 2017): e.g. at $\lambda = 1$, $I_{\Delta} = \mathcal{O}(N^2)$ for $k = 1$, $k = \lfloor \frac{N}{2} \rfloor$, or $k = N$, and $I_{\Delta} = \mathcal{O}(N^4)$ for $k = \lfloor \frac{N}{4} \rfloor$ or $k = \lfloor \frac{3N}{4} \rfloor$.

This proves that the anisotropy at $|\Delta| = 1$ can be precisely estimated also measuring only a single spin of the chain, with the optimal sensitivity attained by a measurement of the single spin magnetization along the z axis. Also the QFI of two contiguous spins at arbitrary dissipation strength λ is expected to scale superextensively at the critical points $|\Delta| = 1$, because the derivative of the current $i\langle \sigma_k^+ \sigma_{k+1}^- - \sigma_k^- \sigma_{k+1}^+ \rangle$ shows a divergence when approaching $|\Delta| = 1$ from the phase $|\Delta| > 1$ (Prosen, 2015). Thus, a precision estimation of the anisotropy at $|\Delta| = 1$ can be inferred by measuring the above current, and the two-contiguous-spin QFI is compatible with a quadratic scaling in N from preliminary numerical computations at $\lambda = 1$ (Marzolino and Prosen, 2017).

For either small λ or small μ the steady states ρ_{∞} of the above models are perturbations of the completely mixed state and thus not entangled (Bengtsson and Życzkowski, 2006). In the XX model $\gamma = \Delta = 0$, there is no nearest neighbor spin entanglement for a wide range of parameters (Žnidarič, 2012).

NESS probes are favorable because QFI is superextensive in a whole phase and not only at exceptional parameters. Moreover, the distinguishability of NESS via Fisher information, thus detectability of non-equilibrium criticality and metrological performances, are enhanced compared to thermal equilibrium systems: e.g. the BKT QPT at $\Delta = \frac{1}{2}$ in the ground state of the XXZ model does not correspond to superextensive QFI (Chen *et al.*, 2008; Sun *et al.*, 2015).

D. Adaptive measurements

Since critical points without critical phases are isolated values, they should be known in advance in order to set the system at criticality and benefit of superextensive QFI. A partial solution is an adaptive approach (Mehboudi *et al.*, 2016) that we generalize to PTs under reasonable conditions. The idea is to perform several estimates changing the thermal state or the NESS, in particular the critical point, at each step according to previous estimates, in order to approach the PT ensuring enhanced sensitivity of the control parameter estimation.

Consider any PT with control parameter θ to be estimated, and assume the QFI $I_{\theta-\theta_c} = \frac{\xi N}{|\theta-\theta_c|^\alpha}$ close to a critical point θ_c with $\alpha > 0$ and a prefactor ξ . We use the notation $I_{\theta-\theta_c}$ to make clear that the QFI depends on the difference $\theta - \theta_c$, rather than on θ alone. Sub-shot-noise sensitivity is shown only at $\theta = \theta_c$. Assume also that θ is initially known within a fixed interval $\theta \in [\theta_{\min}, \theta_{\max}]$ enclosing the value θ_c which can be controlled by other system parameters. First, set the critical point to $\theta_c^{(1)} = \theta_{\max}$, ensuring that $\theta < \theta_c^{(1)}$, i.e. one is on a well-defined side of the phase transition. Find a first estimate $\theta_{\text{est}}^{(1)}$ for the parameter θ , with an uncertainty that saturates the quantum Cramér-Rao bound. Use therefore a number of measurements M that ensures a small error compared to the original confidence interval, $\sigma(\theta_{\text{est}}^{(1)}) = \left(\frac{1}{MI_{\theta-\theta_c^{(1)}}} \right)^{1/2} \ll (\theta_{\max} - \theta_{\min})$, where once more $\sigma(\theta_{\text{est}}^{(1)}) = \text{Var}(\theta_{\text{est}}^{(1)})^{1/2}$ is the standard deviation of the estimate. Then, update the critical parameter to $\theta_c^{(2)} = \theta_{\text{est}}^{(1)} + \sigma(\theta_{\text{est}}^{(1)})$. Since $\sigma(\theta_{\text{est}}^{(1)}) \ll (\theta_{\max} - \theta_{\min})$, the new critical point $\theta_c^{(2)}$ is now much closer to the true value of θ than $\theta_c^{(1)}$, assuming that the obtained estimate $\theta_{\text{est}}^{(1)}$ (which is random and only on average agrees for an unbiased estimator with θ) is indeed within an interval of order $\sigma(\theta_{\text{est}}^{(1)})$ of the true θ . Hence, in the next round, the QFI should be substantially larger. Perform then again sufficiently many times a POVM that allows saturating the quantum Cramér-Rao bound for a new estimate $\theta_{\text{est}}^{(2)}$, $\sigma^2(\theta_{\text{est}}^{(2)}) = \frac{1}{MI_{\theta-\theta_c^{(2)}}} \simeq \left(\frac{\sigma(\theta_{\text{est}}^{(1)})^\alpha}{\xi MN} \right) \propto (MN)^{-\frac{2+\alpha}{2}}$. After k iterations, the sensitivity of the estimate $\theta^{(k)}$ saturating

the quantum Cramér-Rao bound is

$$\sigma^2(\theta_{\text{est}}^{(k)}) = \frac{1}{MI_{\theta-\theta_c^{(k)}}} \propto (MN)^{-\sum_{j=0}^{k-1} \frac{\alpha^j}{2^j}} = (MN)^{\frac{1-(\alpha/2)^k}{\alpha/2-1}}, \quad (91)$$

achieving sub-shot-noise for $\alpha < 2$ with the limiting scaling $\sigma^2(\theta_{\text{est}}) \propto (MN)^{\frac{2}{\alpha-2}}$ for $k \rightarrow \infty$.

Within the above adaptive scheme, control of Hamiltonian parameters, estimations at each step, and re-materialization or re-stabilization of states at each step are assumed to be efficiently implementable. This adaptive measurement was proposed to estimate the critical magnetic field $h = J$ of the XX model at small non-zero temperature when the QPT is smoothened to a phase crossover (Mehboudi *et al.*, 2016), achieving the limiting scaling $\sigma^2(\theta_{\text{est}}) \propto (MN)^{-\frac{4}{3}}$. The optimal estimate is derived from measurements of the magnetization along the z direction $\sum_j \sigma_j^z$, since the magnetic field commutes with the spin interaction restoring the classical picture of Lagrange multipliers and quantities fixed on average. The same estimate and that derived from measurements of the variance of $\sum_j \sigma_j^x$ are nearly optimal for the XY model. It is remarkable that, when the adaptive measurement is applied to the model (74) approaching the critical lines in (76) from the non-critical region, the critical scalings are consistently recovered.

VI. OUTLOOK

Most work on quantum enhanced measurements has investigated the benefits of using quantum entanglement (see (Giovannetti *et al.*, 2011; Paris, 2009; Pezzé and Smerzi, 2014; Pezzè *et al.*, 2016; Tóth and Apellaniz, 2014) for recent reviews). Indeed, under certain restrictive assumptions (see Introduction), entanglement can be shown to be necessary if one wants to improve over classical sensitivity. However, going beyond these restrictive assumptions opens up a host of new possibilities of which we have explored a large number in this review. Given the difficulty of producing and maintaining entangled states of a large number of subsystems, some of these may open up new roads to better sensitivities than classically possible with a comparable number of resources in actual experiments.

We conclude this review by challenging yet another common mind-set in the field (which we could not quite escape in this review either), namely the hunt for faster *scaling* of the sensitivity, in particular the quest for a scaling faster than $1/\sqrt{N}$ with the number of subsystems N , and the goal of reaching “Heisenberg-limited” scaling $1/N$: It should be clear (see also Sec.IV.D and the linear/nonlinear comparison in (Napolitano *et al.*, 2011)) that *scaling* of the sensitivity is not *per se* a desideratum. Any given instrument or measurement is judged by its

sensitivity, not the scaling thereof. When the sensitivity is $\sigma(\theta_{\text{est}}) = \alpha N^d$, it is sometimes argued that the pre-factor α is irrelevant, because a more rapid scaling necessarily leads to better sensitivity for sufficiently large N . While mathematically impeccable, this argument assumes that the scaling persists to sufficiently large N where the possibly small prefactor can be compensated — an assumption that may not be valid in practice. Typically, at some point the model breaks down, and systematic errors arise that scale with a positive power of N and at some point become comparable to the stochastic error quantified by the quantum Cramér-Rao bound. And finally, some large- N catastrophe destroys the instrument and its measurement capability. Such concerns are of course very relevant for real-life experiments, where material properties have to be taken into account. But they may also determine fundamental bounds to achievable precision for various physical quantities that are ultimately linked to the fabric of space-time at extremely small length- and time-scales. Such ideas were advanced early on notably by Wigner, who estimated the ultimate achievable precision of atomic clocks: increasing the energy of the used clock states more and more for improving its precision leads ultimately to the formation of a black hole and hence renders reading off the clock impossible (Wigner, 1957). Similar limitations of this kind exist for measurements of lengths (Amelino-Camelia, 1999; Ng and Dam, 2000), and have been recently explored in more detail for the speed of light in vacuum (Braun *et al.*, 2015).

While current technology is still far from probing such extreme conditions, we nevertheless arrive to the conclusion that the importance of scaling is to give the functional form for an *extrapolation* of the sensitivity, a prediction of how well one could measure if one had a given large N . The validity of this extrapolation will be limited by the range over which the scaling persists, an additional datum not described by the scaling nor by the prefactor. Moreover, there is the possibility of an *optimum* N , beyond which the sensitivity worsens (Nichols *et al.*, 2016). In such a case, the interesting questions are “is the optimum N_{opt} achievable given available resources” and “what is the actual sensitivity at this optimum?” At such an optimum the local scaling is flat, i.e. the smallest possible uncertainty of an unbiased estimate of the parameter is independent of N . Ironically, our discussion of advantageous scaling leads us to the conclusion that the best scaling may be no scaling at all.

The relevance of the actually achievable smallest uncertainty rather than its scaling with the number of resources makes it particularly important that alternatives to the use of massive entanglement be investigated, as so far the number of subsystems that could be entangled

experimentally has remained relatively small. We hope that the present review will stimulate further research in this direction.

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